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EXPERIMENTAL STUDIES OF EXCITED METALLIC PHASES IN STRESSED GERMANIUM

John Elson Furneaux
(Ph.D. thesis)

September 1979

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EXPERIMENTAL STUDIES OF EXCITED METALLIC PHASES
IN STRESSED GERMANIUM

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Experimental Studies of Excited Metallic Phases in Stressed Germanium

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ABSTRACT

The strain-confined electron-hole liquid (SCEHL) is a unique form of matter which forms when photoexcited carriers flow into a strain potential well created by suitably inhomogeneously stressing Ge. Various experimental aspects of this SCEHL were studied.

The formation process for the SCEHL was examined by time-resolved Alfvén wave resonance experiments. The principal results of these experiments are 1) the SCEHL forms quite quickly within a few microseconds, so that the buildup and decay processes are independent; 2) the buildup process depends on the laser pump power indicating that the phonon wind probably contributes to the buildup; 3) the buildup time is the time for the photo-produced carriers to move into the strain potential well; 4) the applied magnetic field has very little effect on the buildup.

The distortion of the SCEHL by a static magnetic field was discovered and studied. Simultaneous infrared imaging and Alfvén wave resonance experiments lead to the discovery of this magnetostriction of the SCEHL. Further studies using an infrared vidicon and an image scanning apparatus demonstrated that the magnetostriction is caused by the "recombination current" which is necessary to keep the SCEHL at a constant density even though electron-hole pairs within it recombine. The externally applied magnetic field deflects these currents causing an induced magnetic moment and a magnetic contribution to the energy. The total energy of the SCEHL is then minimized
by a distortion of the liquid drop usually into an oblate spheroidal shape about the magnetic axis. The experimental results agree very well with a detailed magneto-hydrodynamic theory of this effect. The dynamic nature of the effect is demonstrated conclusively by 1) the dependence of the effect on external pumping conditions, and 2) the time behavior of the effect during the decay of the SCEHL.

The equilibrium properties of the SCEHL were studied by making spectroscopic and threshold measurements using the characteristic recombination luminescence. Fitting the spectroscopic lineshapes for the SCEHL and associated free excitons at temperature T = 3.3 K and stress σ_m ≈ -6 kgf lead to the determination of the SCEHL 1) equilibrium density n_0 = (4.5 ± 0.3) × 10^{16} cm^{-3}, and 2) condensation energy φ = 10.8 ± 1 K. The threshold measurements between 3.3 K and 4.1 K give an independent measurement of φ = 9.2 ± 1 K, in agreement with the spectroscopic measurement.

The de Haas-van Alfvén like magneto-oscillatory luminescence from the SCEHL was studied. The period of these oscillations in 1/H gave an independent measurement of the SCEHL equilibrium density n_0 = (5.2 ± 0.5) × 10^{16} cm^{-3} at T = 1.6 K and σ_m ≈ -6 kgf/mm^2, in agreement with the spectroscopic measurements. Studying the angle dependence of these oscillations confirmed that only one conduction band valley is occupied in the SCEHL. The time dependence of these oscillations provided information about the recombination processes in the SCEHL. It was discovered that a density independent process, probably due to deep traps, is important for the SCEHL.

The metal-insulator transition in highly excited uniformly strained Ge was studied: a spectrally resolved, temperature independent luminescence line at energies between the free exciton and electron-
hole liquid was observed, suggesting the existence of a sharp metal-insulator transition in this material.
CHAPTER I

Introduction
The electron-hole liquid (EHL) is a unique state of matter which forms in photoexcited semiconductors at cryogenic temperatures. As is shown schematically in Fig. I-1, when photons with energies greater than the band gap are absorbed in Ge at low temperatures, electron-hole pairs are formed. These pairs, as a consequence of their opposite charge, bind rapidly into positronium-like species called free excitons (FE). As was first proposed by Keldysh,\textsuperscript{1.1} if the density of excitons is made large enough by exciting the Ge with a sufficient photon flux, and if the temperature is low enough, the FE will condense into a cloud of small (1 - 10 μm) liquid electron-hole drops (EHD) in much the same way as gases such as water vapor condense into liquid fog droplets. This electron-hole liquid (EHL), however, has many novel properties.\textsuperscript{1.2-1.5} Unlike most classical liquids it is a cold compensated metallic plasma with a pair density of about $2 \times 10^{17}$ cm$^{-3}$, and it has a finite lifetime of about 40 μs.

The most fruitful experimental techniques for studying this liquid are luminescence studies involving the characteristic radiation emitted from these excited species (FE and EHL) during their decay.\textsuperscript{1.2-1.4} The energetics of these processes are shown schematically in Fig. I-2. The indirect gap in Ge is 744 meV at liquid helium temperatures. In order to conserve crystal wave vector a band edge phonon, in this case a longitudinal acoustic (LA) phonon of 27 meV, must be emitted along with the photon which is characteristic of the appropriate excited state. The FE is bound with respect to the free pairs (FP) at the unperturbed single particle gap by $E_x \approx 4$ meV. The EHL in unstressed Ge is characterized by a condensation energy, $\phi \approx 2$ meV, between the exciton and the Fermi
level $E_{F\perp}$ of the EHL and a Fermi energy

$$E_F = E_F^e + E_F^h \approx 6 \text{ meV}, \quad (\text{I-0})$$

where $E_F^e$ and $E_F^h$ are the Fermi energies of the electrons and holes respectively in the degenerate Fermi fluid, the EHL. Thus, recombination luminescence for FE is observed at 714 meV and for the EHL at 709 meV.

The changes in band structure caused by a uniaxial stress in Ge have a profound effect on these excited species. The lowest conduction-band ellipsoids are shown schematically in Fig. I-3. There are four distinct degenerate ellipsoids oriented along the crystalline (111) directions. The maximum in the valence-band occurs at $k = 0$ where the heavy and light hole bands are degenerate. The effect of uniaxial stress on these bands is shown schematically in Fig. I-4. For instance, a uniaxial stress along one of the crystalline (111) directions will lower the corresponding ellipsoid in energy. Any uniaxial stress will break the degeneracy in the valence-band. In moderate stress ($\sim 10$ kgf) the conduction and valence bands are split by more than the respective electron and hole Fermi energies in the EHL, thus decreasing the number of populated bands.

Because the kinetic, exchange, and correlation energies of the EH pairs within the EHL depend sensitively on the degeneracy of the bands, it is expected that the binding energy, equilibrium density, and other parameters will strongly depend on strain. For the Ge[4:2], Ge[1:2], and Ge[1:1] configurations (see Fig. I-4), Vashishta et al. I.6,I.7 have calculated equilibrium densities of $2.2 \times 10^{17}$, $0.69 \times 10^{17}$, and $0.11 \times 10^{17}$ cm$^{-3}$, respectively. The density has been greatly reduced, and the EHL is only slightly bound for Ge[1:1]. The reduction in binding with density occurs mainly because the kinetic energy term is much larger for bands which are not
Kelso has extended these calculations to intermediate densities corresponding to a continuous variation of the stress, and there is satisfactory agreement with experiment.

Since these strains correspond to a lowering of the energy of an EH pair in any of the states FP, FE, or EHL these pairs will be attracted to regions of the crystal where there is a maximum strain. I use the term FP to refer to a free electron and a free hole without any electrostatic interaction. This leads to motion of EHDs in a strain gradient which was first reported by Alekseev et al. Therefore, by suitably straining the crystal it is possible to collect relatively large amounts into a single drop, producing the strain-confined electron-hole liquid (SCEHL).

To describe the shift of the band edge with strain in Ge, the usual second rank strain tensor with a typical element

\[ \varepsilon_{xy} = \frac{1}{2} \left( \frac{\partial u_x}{\partial y} + \frac{\partial u_y}{\partial x} \right) \]  

is defined where \( u_x \) is the local displacement in the x direction. In the case of the inhomogeneous strains which are of interest here, the elements of this tensor are functions of position within the crystal. The local stress tensor \( \sigma \) is related to \( \varepsilon \) through the stiffness tensor \( C \):

\[ \sigma = \varepsilon 

The deformation potentials \( \Xi_u \) and \( \Xi_d \) introduced by Herring and Vogt are used to describe the shift of the \( i \)th conduction band with strain:

\[ \Delta E_c^i = \Xi_u \hat{a}_i \cdot \varepsilon \cdot \hat{a}_i + \Xi_d \text{Tr} \varepsilon, \]

where \( \hat{a}_i \) is the unit vector pointing along the \( i \)th valley (i.e., a \( \langle 111 \rangle \) direction), and \( \text{Tr} \varepsilon = \varepsilon_{xx} + \varepsilon_{yy} + \varepsilon_{zz} \). This may be written in the form
\[
\Delta E^i_c = \Delta E^\text{mean}_c + \varepsilon_o^i \varepsilon^o,
\]  
(I-4)

where

\[
\Delta E^\text{mean}_c = (\varepsilon_d + \frac{1}{3} \varepsilon_u) \text{Tr} \varepsilon^o
\]  
(I-5)

is the mean shift of all four conduction bands, and relative shift of the \(i\)th Valley, \(\varepsilon^i_o\), is defined by

\[
\varepsilon^i_o = \hat{a}_i \cdot \varepsilon \cdot \hat{a}_i - \frac{1}{3} \text{Tr} \varepsilon.
\]

These shear strains have the dominant influence on the band gap shift. The deformation potentials \(a\), \(b\), and \(d\) defined by Picus and Bir\(^{1,15}\) are used to describe the shift in the \(j\)th valence-band edge:\(^{1,13}\)

\[
\Delta E^j_v = a \text{Tr} \varepsilon + (-1)^j |E_{\varepsilon \varepsilon}|,
\]
(I-6)

\[
E_{\varepsilon \varepsilon}^2 = 1_4 b^2 [(\varepsilon_{11} - \varepsilon_{22})^2 + \text{c.p.}] + d^2 (\varepsilon_{12}^2 + \text{c.p.})
\]  
(I.7)

where c.p. stands for cyclic permutations with respect to the indices 1, 2, and 3 which stand for the \((100)\) axis of the crystal. Values for these deformation potentials taken from Balslev\(^{1,10}\) are

\[
\begin{align*}
\varepsilon_u &= 16.2 \pm 0.4 \text{ eV} \\
b &= -1.8 \pm 0.3 \text{ eV} \\
d &= -3.7 \pm 0.4 \text{ eV}
\end{align*}
\]  
(I-8)

\[
\varepsilon_d + \frac{1}{3} \varepsilon_u - a = -2 \pm 0.5 \text{ eV}.
\]

For an arbitrary strain, these strain-split band edges give rise to a number of band gaps: \(E^{ij}_g = E_g(\varepsilon=0) + \Delta E^i_c - \Delta E^j_v\). Carriers will relax to the minimum gap:\(^{1,13}\)

\[
E_g(\varepsilon) = E_g(0) + (\varepsilon_d + \frac{1}{3} \varepsilon_u - a) \text{Tr} \varepsilon - |E_{\varepsilon \varepsilon}| - \varepsilon_u [\max_1(-\varepsilon^i_o)].
\]  
(I-9)

The resulting shift in the minimum band gap \(\Delta E_g\) is plotted versus uniaxial stress for the three principal crystal directions in Fig. I-5.
From the magnitudes of the deformation potentials in Eq. (I-8), it is obvious that the last term in Eq. (I-9) due to the conduction band splitting is the main source of the band-gap shift. The octahedral shear strain is defined as

\[ \varepsilon^m = \max_i (-\varepsilon_i) , \]  

which is greater than zero because compressional strains are taken as negative. Therefore in a nonuniform strain the free carrier energy will be a function of position and will minimize at the point of maximum octahedral shear strain which also corresponds to the point of maximum octahedral shear stress.

The shift of the energy of the EHL has been measured in uniformly strained Ge by a number of groups. The data of Benoit et al. are plotted with the model of Markiewicz et al. in Fig. I-6. It can be seen that the model explains the experimental results quite well. More sophisticated theory provides an even better explanation. For low stresses the EHL energy does not follow the band gap because the binding energy is decreasing due to the loss of degeneracy. Once the strain-split bands are completely depopulated the EHL energy shift (\( \Delta E_{\text{EHL}} \)) is approximately parallel to the band-gap energy shift. Balslev has also found that the FE energy shift is parallel to the band gap energy shift at these stresses. Thus above a certain threshold stress (\( \sim 3 \text{ kgf/mm} \)), excitons and EHL will be attracted to regions of maximum \( \varepsilon^m \).

The most convenient method for creating a strain potential well corresponding to a region of maximum \( \varepsilon^m \) inside the crystal has been to apply a Hertzian contact stress. This is a contact stress between two ellipsoidal elastic solids, in our case the Ge sample and a rounded nylon plunger as pictured in Fig. I-7. This provides a
maximum $\varepsilon_0$ inside the crystal. Markiewicz\textsuperscript{1.13} has made detailed calculations using a sophisticated finite-element analysis computer program. This gives the results pictured in Fig. I-8. Both $\Delta E_g$ and $\Delta E_{\text{EHL}}$ are shown as a function of position within the crystal. An approximately parabolic strain-potential well is formed inside the crystal.\textsuperscript{1.13} (Equal energy contours obtained by Markiewicz are shown in Fig. III-1.) These theoretical predictions for the strain well are verified extensively by Markiewicz et al.\textsuperscript{1.13} using infrared birefringence studies and EHL luminescence imaging studies.

The strain-confined electron-hole liquid was discovered by observing Alfvén wave resonance absorption of microwaves in non-uniformly stressed Ge.\textsuperscript{I.19} An Alfvén wave is an electromagnetic wave which propagates with wavelength $\lambda_a$ in a compensated plasma in the presence of a static magnetic field. While the full theory of these resonances is quite complicated,\textsuperscript{I.20,I.21} the basic ideas are easily understood. In order to excite the fundamental standing Alfvén wave in a spherical drop of EHL of radius $R$, the Alfvén wavelength must approximately match the drop radius:

$$\lambda_a \approx R. \quad (I-11)$$

Since the wavelength $\lambda_a$ within the plasma sphere is related to the free space wavelength $\lambda_0$ through the dielectric constant of the EHL $\varepsilon_{\text{EHL}}$

$$\lambda = \lambda_0 / \sqrt{\varepsilon_{\text{EHL}}}, \quad (I-12)$$

the frequency and magnetic field-dependent $\varepsilon_{\text{EHL}}$ must be calculated. This simple Alfvén theory is discussed below.

For a metal with a single mobile carrier with a density $n$, mass $m_e$, and charge $q = -e$, the dielectric constant for circularly polarized
waves is \textsuperscript{1.20,1.21}

\[
\varepsilon = 1 - \frac{\omega_p^2}{\omega(\omega - \omega_c)}
\]  
(I-13)

where the plasma frequency is defined by

\[
\omega_p^2 = \frac{4\pi n e^2}{m_e}
\]  
(I-14)

and the cyclotron frequency by

\[
\omega_c = \frac{qH}{mc}
\]  
(I-15)

where \(H\) is the magnetic field. However, for the case of the EHL there are also holes with mass \(m_h\) and charge \(e\). This leads to

\[
\varepsilon_{EHL} = 1 - \frac{\omega_{pe}^2}{\omega(\omega - \omega_{ce})} - \frac{\omega_{ph}^2}{\omega(\omega - \omega_{ch})}.
\]  
(I-16)

In the experiments discussed here, \(\omega_p \gg \omega_{ci} \gg \omega\) and lowest order in \(\omega/\omega_{ci}, i = e\) or \(h\), \(\varepsilon_{EHL}\) becomes the Alfvén dielectric constant

\[
\varepsilon_a = \frac{\omega_{pe}^2}{\omega_{ce}^2} + \frac{\omega_{ph}^2}{\omega_{ch}^2} \approx \frac{n}{H^2}.
\]  
(I-17)

Therefore by Eqs. (I.11), (I.12), and (I.17)

\[
R \propto \lambda_a = \frac{\lambda_0}{\varepsilon_a^2} \propto \lambda_0 H/n^{1/2}
\]  
(I-18)

or

\[
H \approx n^{1/2} R_\omega.
\]  
(I.19)

The resonance position in magnetic field for fixed \(n, \omega,\) and shape is a good indication of the drop size. A more exact theory \textsuperscript{1.20,1.21} yields for the relation between the resonant field and the radius of the compensated plasma sphere at a density of \(n = 0.5 \times 10^{17}\) and frequency \(2\pi\omega = 25\) GHz

\[
R(\mu\text{m}) = 14[H(\text{kOe}) - 0.3].
\]  
(I.20)
In order to gain more information about this novel state of matter, the electron-hole liquid, a number of experiments were performed on ultrapure Ge samples prepared from single crystals grown by W. Hansen and E.E. Haller and kindly furnished us. Table I-1 is a list of the characteristics of the particular samples used in the experiments presented here. Most of the experiments were done on the sample CR36.

The mechanism of the SCEHL formation and growth is an obvious experimental question. We studied the kinetics of this growth by the use of Alfvén wave resonance experiments and found that the SCEHL forms much more quickly than it decays. Formation times are on the order of μs, and decay times are on the order of ms. These experiments are presented in detail in Chapter II.

Infrared imaging experiments in a magnetic field, while attempting to measure the equilibrium density of the SCEHL with Alfvén waves, lead to the discovery of the marked distortion of the EHL by a static magnetic field. A complete magneto-hydrodynamic theory of this magnetostriction was worked out by Markiewicz.22,23 We performed a number of experiments on the SCEHL which turn out to agree with this theory. A review of this theory and our complete experimental results are presented in Chapter III.

In order to more thoroughly understand the equilibrium properties of the SCEHL a number of luminescence experiments were performed. First, spectroscopic measurements and lineshape analysis lead to the determination of the equilibrium density and condensation energy of the SCEHL. The FE lineshape predicted for the strain-confined free exciton gas was also confirmed. Second threshold studies of the integrated FE and EHL luminescence from the strain potential well confirmed the spectroscopic determination of the liquid condensation energy.
These experiments are presented in Chapter IV.

Magneto-oscillatory luminescence experiments were performed on the SCEHL because they had proven so fruitful for the unstressed EHL. They proved equally fruitful for us. By studying the period of these de Haas-van Alfvén like oscillations, the electron Fermi energy was found. This allowed an independent determination of the equilibrium density of the SCEHL. The angle dependence of these oscillations confirmed that only one conduction band valley is occupied in the SCEHL. Finally, the time dependence of these oscillations provided information about the decay processes of the SCEHL. These experiments are discussed in Chapter V.

Ge uniaxially strained along a (111) direction provides an almost ideal media for studying metal-insulator transitions in highly excited semiconductors. The valence band degeneracy has been lifted and the FE can therefore be handled much more easily theoretically. The experimental situation is also simplified; moderate excitation levels and temperatures below 4.2 K are sufficient to see effects. I. Balslev and I therefore performed experiments on (111) uniformly stressed Ge. Some theory on the nature of the metal-insulator transition in Ge[1:1] and our results are presented in Chapter VI.

Chapter VII contains my conclusions.
### Table I-1

<table>
<thead>
<tr>
<th>Sample No.</th>
<th>Dimensions (mm)</th>
<th>Face</th>
<th>Stress</th>
<th>Boule No.</th>
<th>Special Preparation</th>
</tr>
</thead>
<tbody>
<tr>
<td>CR 14</td>
<td>3.95 disc x 1.40</td>
<td>(100)</td>
<td>(011)</td>
<td>145</td>
<td></td>
</tr>
<tr>
<td>CR 36</td>
<td>4 x 4 x 2</td>
<td>(110)</td>
<td>(111)</td>
<td>145</td>
<td></td>
</tr>
<tr>
<td>CR 37</td>
<td>4 x 4 x 2</td>
<td>(110)</td>
<td>(111)</td>
<td>145</td>
<td></td>
</tr>
<tr>
<td>CR 38</td>
<td>3.85 x 3.85 x 1.75</td>
<td>(110)</td>
<td>(111)</td>
<td>146</td>
<td>Syton polish</td>
</tr>
<tr>
<td>CR 50</td>
<td>3.85 x 3.95 x 2.80</td>
<td>(110)</td>
<td>(111)</td>
<td>145</td>
<td></td>
</tr>
<tr>
<td>MI 1</td>
<td>0.7 x 0.7 x 11</td>
<td>(111)</td>
<td></td>
<td>437</td>
<td></td>
</tr>
</tbody>
</table>

**a** All samples were polished and etched in White etch, 3HNO₃:HF.

**b** The boule number is the number designated by W. Hansen and E. Haller at LBL for the characterization of crystals they have grown. The characterizations given below are from E.E. Haller by private communication.

**c** Boules 145 and 146 are virtually identical dislocation free (100) grown, net acceptor concentration $N_A \approx 2 \times 10^{11}$ probably due to a divacancy plus hydrogen.

**d** Boule 437 is lightly dislocated, $\sim 100 - 1000$ etch pits/cm²; net donor concentration $N_D \approx 10^{11}$ probably due to phosphorus; silicon impurities $\sim 10^{16}$. 
CHAPTER I: References


I.5 T.M. Rice, ibid., p. 1.


I.18 H. Hertz, J. Math. (Creille's J.) 92, 156 (1881).
CHAPTER I: Figure Captions

Fig. I-1 Model of the processes inside a Ge crystal which lead to the formation of electron-hole drops. The existence of biexcitons is not yet confirmed in Ge.

Fig. I-2 Band structure $E(k)$ for Ge with schematic representations of the free exciton energy $E_x$, and the electron-hole drop Fermi level at $E_B$, with total Fermi energy $E_F$ and condensation energy $\phi$. The band-edge longitudinal acoustic phonon energy is $E_{LA}$.

Fig. I-3 Lowest conduction-band ellipsoids in Ge, showing the constant energy surfaces $E(k)$ orientated along the (111) crystal axes in the first Brillouin zone; the ellipsoids are shown cut at the zone boundary. Edges of the cube are shown parallel to (100) crystal directions. Lattice constant $a = 5.65\, \text{Å}$ for Ge (from Ref. I.13).

Fig. I-4 Splitting of the electron and the hole bands in Ge under uniaxial compression. For example, stress along a (111) direction lowers one electron ellipsoid in energy relative to the other three. For a medium stress this splitting exceeds the Fermi energy of the electrons, and only one ellipsoid is populated; since both hole bands are still populated, we denote this configuration as Ge[1:2]. At higher stresses only one hole band is populated, denoted by Ge[1:1]. Unstressed Ge is denoted by Ge[4:2].

Fig. I-5 Shift in band-gap energy with uniaxial stress at 80 K in Ge: $\Delta E_g = E_g(\varepsilon) - E_g(0)$ from Eq. (I-9). Lines drawn have slopes corresponding to the data of Ref. I.10 (from Ref. I.13).
1.4 mm as a function of Z. (b) Potential wells at \( z = 0.5 \text{ mm} \) as a function of \( y \). The slight asymmetry in \( y \) is because Ge is not symmetric under reflection in a (112) plane. The solid lines are parabolas, showing that the well can be readily approximated by a spherical potential \( V \propto r^2 \) where \( r \) is the radial distance from the minimum point, the point of maximum octahedral shear stress.

**Fig. I-6** Shift of EHL energy with uniaxial stress: \( \Delta E_{EHL} = E_{EHL}(\varepsilon) - E_{EHL}(0) \). Data points are taken from Ref. I.16. Solid lines are from Ref. I.13 (theory).

**Fig. I-7** Experimental arrangements for applying inhomogeneous stress to Ge samples. (a) and (b) "permanent stress" holder; (c) variable stress apparatus. For details, see Ref. I.13.

**Fig. I-8** (a) Insert: nylon plunger pressing Ge crystal, producing a strain potential minimum inside the crystal at \( z = 0.5 \text{ mm} \); the curves show the band-gap shift, \( E_g \), and EHL energy shift \( \Delta E_{EHL} \) calculated by Markiewicz I.13 for \( F = 10 \text{ kgf} \) and \( 2a = 1.4 \text{ mm} \). (b) Potential wells at \( z = 0.5 \text{ mm} \) as a function of \( y \). The slight asymmetry in \( y \) is because Ge is not symmetric under reflection in a (112) plane. The solid lines are parabolas, showing that the well can be readily approximated by a spherical potential \( V \propto r^2 \) where \( r \) is the radial distance from the minimum point, the point of maximum octahedral shear stress.
Fig. 1-1

- 714 meV LIGHT
- DROP EMBRYO
- EXCITING LIGHT
- BIEXCITONS (?)
- 709 meV LIGHT
- MACROSCOPIC ELECTRON-HOLE DROP (EHD) 5 MICRONS
- HOLE ELECTRON FREE EXCITON (FE)
GERMANIUM

$E_g = 744 \text{ meV}$

$\phi \approx 2 \text{ meV}$

$E_F \approx 6 \text{ meV}$

$E_B$:

$E_X \approx 4 \text{ meV}$

$E_{LA} = 27 \text{ meV} \text{ phonon}$

Fig. I-2
STRESS DIRECTION

\[
\begin{array}{ccc}
\langle 100 \rangle & \langle 110 \rangle & \langle 111 \rangle \\
Ge (4:2) & Ge (2:2) & Ge (1:2) \\
\end{array}
\]

MEDIUM STRESS
\(\sim 3 \text{ kgf/mm}^2\)

\[
\begin{array}{ccc}
\text{Ge (4:1)} & \text{Ge (2:1)} & \text{Ge (1:1)} \\
\end{array}
\]

HIGH STRESS
\(\sim 10 \text{ kgf/mm}^2\)

\[
\begin{array}{ccc}
\text{XBL 796:6467} \\
\end{array}
\]

Fig. I-4
Fig. 1-5
Fig. I-6
CHAPTER II

Buildup of the Strain-Confined Electron-Hole Liquid
Once it had been established that the strain-confined EHL existed, we decided to investigate its formation process. The most appropriate experimental technique for initial investigations was Alfvén wave resonance studies. This technique can very effectively probe quite small volumes of SCHEL and has good time resolution (~0.1 μs) due to the inherent speed of the microwave receiver. I therefore set up the experimental apparatus pictured in Fig. II-1. A model LD-167 single heterojunction stacked diode laser of wavelength 0.904 μm was pulsed at 500 Hz with a pulse width of 0.1 μs and peak energy of 7.2 ergs per pulse absorbed into the crystal. The laser was mounted on a precision x-y translation stage so the dependence of the buildup on laser pump position could be determined. A quartz lens was mounted on the tail of the dewar, and the laser was positioned under the dewar such that the 0.4mm × 0.4mm laser surface was focused onto the germanium sample surface with a magnification of 1. This calibrated the movement of the laser with the x-y translator to correspond to the movement of the pumping spot on the sample.

Typical microwave absorption spectra are shown in Fig. II-2 for various times after the light pulse. The drop buildup and decay are measured by the growth and decay of the principal resonance peak, $H_c$. Figure II-3 is a plot of field vs. time for the principal resonance of Fig. II-2. It can be seen that the Alfvén resonance position increases rapidly to a maximum field $H_{max}$, corresponding to a maximum drop size, and then decays exponentially with a characteristic 1400 μsec time constant. If we assign the principal resonance at $H_c$ in Fig. II-2 as the lowest magnetic dipole mode of a sphere, the relation between radius and field, as was shown in Chapter I is
\[ R(\mu m) = 14[H_c(kOe) - 0.3] \]  
(II-1)

for a liquid density of \(0.5 \times 10^{17} \text{ cm}^{-3} \). For this mode assignment the maximum resonant field of 13 kOe corresponds to a drop radius of 178 \(\mu m\). This is approximately the same size I obtain from assuming 100% pumping efficiency into the strain well. The maximum laser pulse of 7.2 ergs corresponds to about \(10^{12} \) electron-hole pairs and a drop of 170 \(\mu m\) radius. The observed decay time of 1400 \(\mu s\) corresponds very well with observed radius and volume decay times observed in this sample and other similar samples. II.4, II.8, II.9 Typically volume decay times of 400-500 \(\mu s\) are observed corresponding to radius or Alfvén resonance decay times of 1200-1500 \(\mu s\). Figure II-4 shows a plot of the maximum resonant field \(H_{\text{max}}\) as a function of absorbed laser energy \(P_a\). As would be expected for this mode assignment, \(H_{\text{max}} \propto P_a^{1/3}\). The deviation shown at low power will be discussed later.

Figure II-5 shows the buildup for various energies per pulse absorbed. It can be seen that for a lower power it always takes longer to attain a given drop size. If the buildup process were independent of laser power, then the resonance position as a function of laser power and time for the buildup would be

\[ H_c = H_{\text{max}} f(t_d) \]  
(II-2)

where \(f(t_d)\) is a function only of the delay time, \(t_d\). This hypothesis is tested in Fig. II.6. The relative size, \(H_c/H_{\text{max}}\) from Fig. II-5 is plotted vs. time. Because the five curves from Fig. II-5 remain distinct when replotted in this manner and do not fall on some universal curve \(f(t_d)\), I conclude that the buildup process is also dependent on the excitation energy. The data in Fig. II-6 can be explained by a pulse energy independent delay time of about 1 \(\mu s\) and then an energy dependent rise to the maximum drop size in the next 3-5 \(\mu s\).
For higher pulse energies there is a faster energy dependent build-up. The data displayed in Fig. II-7 tests this hypothesis. Energy per pulse is plotted versus delay time for various resonant fields. Then the initial delay time for the first carriers to get into the well and form an infinitesimally small drop, $t_d(H_c=0)$, was extrapolated from the data as shown in Fig. II-8. The data are replotted as absorbed energy per pulse versus resonant field for various times. At the longest time, 5 ms, it was found that $H_{\text{max}} \propto P_a^{1/3}$, Fig. II-4. For shorter times the data are consistent with

$$H_{\text{max}} \propto P_a^{1/3} = H_c + H_0(t_d) \quad (\text{II-3})$$

where $H_0(t_d)$ is the time dependent offset field. The above conclusion is tested in Fig. II-9 where $H_{\text{max}} - H_c$ is plotted as a function of delay time for the data from Fig. II-5. The surprising result is found that within experimental error and inaccuracies in finding $H_{\text{max}}$, $H_{\text{max}} - H_c$ is independent of laser power. This means that the drop size

$$R \approx H_c \approx H_{\text{max}}(P_a) - H_0(t_d). \quad (\text{II-4})$$

Thus the extrapolation curves to $H_c=0$ shown in Fig. II-8 are justified. These extrapolations are plotted in Fig. II-7 (error bars forming the far lefthand curve), and a smooth curve is drawn through them. This is consistent with an approximately energy-independent initial delay time for all but the lowest pump powers; however, other extrapolations are possible.

This delay time has been explained by Manoliu and Kittel as the time for EHL formed at the surface by the exciting pulse to move into the strain well. This time is long because a drop strongly emits phonons due to Cerenkov radiation at drop velocities much below the velocity of sound. This is because the constituent electrons and
holes within the liquid are moving at velocities up to the Fermi velocity and therefore are traveling faster than the speed of sound even though the drop as a whole is traveling slower. This damping mechanism becomes progressively stronger as the drop is accelerated and adequately explains the $\sim 1 \mu s$ delay as being the time for the EHL to move the $\sim 1 \text{ mm}$ from the surface into the strain well.

The further delay for the drop to grow to a finite size in the strain well is probably due to the initial distribution of small drops formed by the $0.4\text{mm} \times 0.4\text{mm}$ laser spot on the sample surface. The buildup is most likely faster for higher laser powers due to the phonon wind from the initial distribution and laser pulse.

The time of flight mechanism for the buildup can be tested by studying the buildup of the strain confined drop as a function of the position of the laser spot on the sample. Typical results for maximum drop size are shown in Fig. II-10. Because the buildup occurs on a time scale very short compared to the decay, I will consider the decay to be characterized by

$$H_c(\text{decay}) = H_{\text{max}} \exp \left( -\frac{t_d}{\tau} \right),$$

where $\tau$ is the radius decay time typically $1400 \mu\text{sec}$. This implies that $t_d$ as plotted in Fig. II-10 is equal to $\ln(H_{\text{max}}/H_c) + \text{const.}$, and therefore Fig. II-10 is a reliable measure of $H_{\text{max}}$ vs. position.

The data in Fig. II-10 dramatically show the effect of the strain well on drop size. The changes are approximately symmetric about the point of maximum size for movement in the $y$ direction perpendicular to the strain but are markedly asymmetric for movement in the $z$ direction parallel to the strain. This is the sort of behavior predicted by
Similar data for different energies per pulse are shown in Fig. II-11. A change in the intensity distribution of the laser is the most probable cause for the shift in the position of maximum size for the lowest power. Figure II-12 shows the maximum size ($t_d$ in decay), as a function of absorbed energy per pulse for various average distances of the laser from the position of maximum efficiency from Fig. II-11. It can be seen that the dependence is quite complicated, and there is insufficient data to completely understand it. The maximum size is obviously a function of both position and pulse energy in an inseparable way. Some insight into the problem can be gained from the data in Fig. II-4. The open circles are for pumping at the position of maximum efficiency whereas the triangles are for pumping at a position 0.35 mm away in the $y$ direction. The deviations from $H_{\text{max}} = P^{1/3}$ at lower energies for the displaced spot can be explained by a loss of a certain fixed amount of energy in each pulse.

Figure II-13 shows contours of constant pumping efficiency for two different laser powers. These contours were taken by moving the laser spot on the crystal such that a constant $t_d$ was found in decay for a fixed $H_c$. The $t_d$ for buildup was then recorded at all possible positions. It can be seen that even for a constant final drop size and laser power, the buildup depends markedly on position. For the outer contour variations between 8 and 23 $\mu$s are observed, while for the inner contour even larger variations of between 2.8 and 19 $\mu$s are observed. These buildup times can be explained as a complicated convolution between the initial drop distribution and the inhomogeneous strain field with a velocity dependent damping.
Some further information can be gained by studying contours. In this case the effect of the direction of the magnetic field, if any, on the buildup is found. This is illustrated in Fig. II-14 where contours of constant buildup time for constant $H_c$ and laser energy are plotted for two field orientations, $\hat{H}$ parallel to the stress and $\hat{H}$ perpendicular to the stress. It can be seen that there are very small differences between the contours indicating that the field has very little if any effect on the buildup, even though the buildup time of $30 \mu s$ is quite long. This means that the magneto-hydrodynamic damping proposed by Salvan et al. II.15 does not seem to be particularly important in this case. However, there seems to be some loss of pumping efficiency with magnetic field as seen in Fig. II-2.

The effect of carrier mass anisotropy on the motion of drops in the strain well is shown in Fig. II-15 from Markiewicz. II.16 This shows how the equal energy contours of Fig. II-15a are distorted. Carriers and small drops will flow perpendicular to the contours illustrated in II-15b rather than perpendicular to the contours in Fig. II-15a. This happens because although the force, $\hat{F}$, is perpendicular to the contours in Fig. II-15b, the acceleration

$$\hat{a} = \hat{m}^{-1} \cdot \hat{F}$$  \hspace{1cm} (II-6)

where $\hat{m} = \hat{m}_e + m_n \hat{I}$ and $m_n = 0.3 m_e$. The electron effective mass tensor $\hat{m}_e$ is quite anisotropic for stressed Ge; $m_{\parallel} = 1.58 m_e$ (along the $\langle 111 \rangle$ axis), $m_\perp = 0.082 m_e$. This leads to the distortions in the acceleration potential, $E^*$, pictured in Fig. II-15b. $E^*$ is defined so that $m_{\parallel} \hat{a} = -\hat{\nabla} E^*$. There is a strong tendency for drops to pile up along the $\langle 111 \rangle$ axis and flow into the strain well predominantly along that axis because of the large longitudinal mass. This can be seen
experimentally in Fig. II-16. It is obvious especially from Fig. II-16a and II-16c that drops are channeled as expected from the contours in Fig. II-15b.

Although quantitative information is very difficult to obtain about the buildup process, it is quite well understood qualitatively. Photons from the pumping light are absorbed within about a micron forming a dense fairly hot cloud of electrons and holes. This cools quite quickly ($\sim 10^{-9}$ sec) and forms small EHD. These EHD then move under the combined force of the strain gradient and phonon wind into the strain well at near the velocity of sound ($\sim 10^5$ cm/sec), and are strongly channeled into (111) directions due to the mass anisotropy. Finally, there seems to be very little effect of the magnetic field direction on the buildup as long as it is perpendicular to the direction of the incident light.
CHAPTER II: References


II.7 This relation gives only qualitative size estimates since the drop is markedly non-spherical in a magnetic field as discussed in detail in Chapter III.


II.10 A. Manoliu and C. Kittel, Sol. State Comm. 21, 641 (1977)


II.13 R.S. Markiewicz, Phys. Rev. B (to be published) and references therein.


CHAPTER II: Figure Captions

Fig. II-1 Schematic of experimental apparatus used to measure buildup and decay of EHL with Alfvén wave absorption.

Fig. II-2 Alfvén wave resonances observed at specified times after the 100 nsec pulse. For constant drop density the principal microwave absorption peak occurs at a field $H_c$ approximately proportional to the drop radius. At $t = 0$ approximately $10^{12}$ electron-hole pairs are produced at the crystal surface. After a short delay, the drop radius is seen to increase rapidly to a maximum value, then decay slowly in time. $H$ is in the (100) plane at an angle 25° from (100). $T = 1.65$ K, $v = 25$ GHz and applied stress $\parallel$ (110). Cyclotron resonances of free electrons (E) and holes (H) are prominent initially, but become greatly reduced as electron-hole pairs enter the liquid phase. Sample CR14.

Fig. II-3 Graphs of $H_c$ vs. time after the laser pulse, showing the buildup and decay of the drop radius. The radius decay time of 1400 $\mu$s corresponds to a drop volume decay time of 470 $\mu$s, consistent with the 500 $\mu$s luminescence decay observed for this sample, CR 14.

Fig. II-4 A plot of the maximum drop size as measured by Alfvén wave resonances as a function of absorbed energy per pulse. Stress along (111) direction and $\parallel$ to $\mathbf{H}$; $T = 1.7$ K; $v = 24$ GHz. Circles show laser at position of maximum pumping efficiency. Triangles show results for the laser moved 0.25mm approximately $\perp$ to the stress direction. Sample CR36.
Fig. II-5  Plots of Alfven resonance field, \( H_c \) versus buildup time, \( t_d \), for various laser pulse energies. Same experimental conditions as for the open circles in Fig. II-4. Triangles with points down indicate absorbed energy of 0.7 ergs; squares indicate absorbed energy of 1.4 ergs; triangles with points up indicate an absorbed energy of 2.8 ergs; diamonds indicate an absorbed energy of 5.6 ergs; and open circles indicate an absorbed energy of 7.2 ergs. Sample CR36.

Fig. II-6  Data of Fig. II-5 normalized to \( H_{\text{max}} \) for each energy. It can easily be seen that the buildup is not independent of pulse energy.

Fig. II-7  Plots of energy per pulse versus buildup time for various resonance positions: a) \( H_c = 5kOe \); b) \( H_c = 6kOe \); c) \( H_c = 7kOe \); d) \( H_c = 8kOe \); e) \( H_c = 9kOe \); f) \( H_c = 10kOe \); g) \( H_c = 11kOe \); h) \( H_c = 12kOe \); i) \( H_c = 13kOe \); j) \( H_c = 14kOe \). The solid curve is an extrapolation to \( H_c = 0kOe \); see text for details. All other experimental conditions are the same as in Fig. II-4. Sample CR36.

Fig. II-8  Data of Fig. II-7 replotted as absorbed energy per pulse versus resonant field, \( H_c \), for various buildup delay times \( t_d \). The solid curves are extrapolations of the data to \( H_c = 0 \) assuming \( H_c = H_{\text{max}} - H_0(t_d) \), see text. Symbols are used as follows: open circles 5 \( \mu s \); filled circles 4 \( \mu s \); open square 3 \( \mu s \); filled squares 2.5 \( \mu s \); open triangles points down 2 \( \mu s \); filled triangles points down 1.5 \( \mu s \); open triangles points up 1.25 \( \mu s \), filled triangles points up 1.2 \( \mu s \); open diamonds 1.15 \( \mu s \); filled diamonds 1.1 \( \mu s \), right half-filled circles 1.05 \( \mu s \), bottom half-filled circles 1.0 \( \mu s \), half-filled squares 0.95 \( \mu s \).
Fig. II-9 Data of Fig. II-4 replotted as $H_{\text{max}} - H_c$ versus $t_d$. It can be seen that this is a universal curve approximately independent of pulse energy.

Fig. II-10 Delay time for decay, indicating maximum drop size $t_d = \ln(H_{\text{max}}/H_c) + \text{const.}$ as a function of laser pump position on the crystal. y) laser pump position moved $\perp\text{ to the stress; z)}$ laser pump position moved $\parallel\text{ to the stress. The energy absorbed per pulse is 5.6 ergs and } H_c = 6\text{kOe. All other experimental conditions are as for Fig. II-4. Sample CR36.}$

Fig. II-11 y scans as in the previous figure with $H_c = 6\text{kOe}$ for different absorbed energies: a) 7.2 ergs; b) 3.2 ergs; c) 0.68 ergs. $H \parallel\text{ stress } \langle 110 \rangle$ all other conditions as in Fig. II-2.

Fig. II-12 Data from the previous figure plotted as delay times in decay versus absorbed energy per pulse for different distances from the point of maximum efficiency: a) point of maximum efficiency, b) 0.25mm off in y, c) 0.50mm off in y, d) 0.75mm off in y. It can be seen that there is no simple dependence of drop size on pump power and position.

Fig. II-13 Contours of constant pumping efficiency versus position on the crystal. $H_c = 6\text{kOe}; t_d = 100\text{ sec in delay. a) Absorbed energy is 7.2 ergs per pulse; b) absorbed energy is 1.6 ergs per pulse. The numbers next to the points indicate } t_d \text{ in buildup. The circle and square indicate approximate drop size and laser spot size, respectively. All other experimental conditions as in Fig. II-10.}$

Fig. II-14 Contours of constant buildup time as a function of magnetic field direction: Dots $H \parallel z$, open circles $H \parallel y$. The box in the corner shows the approximate laser spot size. $H_c = 6\text{kOe;}$
\[ t_d = 30 \text{ ms} \quad \text{in buildup; absorbed energy is 7.2 ergs per pulse; } \nu = 24.3 \text{ GHz; } T = 1.35 \text{ K. Very little difference is seen for the different field orientations. Sample CR36.} \]

**Fig. II-15** (a) Calculated electron-hole liquid pair energy in nonuniformly strained Ge. Stress along (111) axis, viewed through (1\(\overline{1}0\)) crystal face. In the unshaded region all electrons are in a single conduction band only. (b) Channeling of carrier flow. Due to effective mass anisotropy in the strain well of Fig. 5(a), the carriers and small drops would actually flow perpendicular to the contour lines drawn here. \( E^* \) is defined so that \( \mathbf{m} \mathbf{a} = -\nabla E^* \), where \( \mathbf{a} \) is the pair acceleration.

**Fig. II-16** Channeling of the carrier flow. The crystal is stressed along a (110) direction and viewed through a (001) face. For enhanced contrast, a negative photographic image is presented so that the glowing drop appears black. The three views correspond to three positions of the laser pumping spot on the sample surface. The arrows show the approximate position of the focused laser spot on the opposite face of the crystal. The carriers appear as a dark cloud of small drops formed about the laser spot. Leading away from the cloud is a narrow channel of drops flowing into the SCEHL near the top of the crystal.
Fig. II-1
a) Buildup

Time (μs)

b) Decay

Time (μs)

Fig. H-3
Fig. II-3
RELATIVE SIZE, $H_c/H_{max}$

BUILDUP TIME, $t_d - \mu s$

Fig. II-6
Fig. II-13
CHAPTER III

Magnetostriiction of the Strain Confined Electron-Hole Liquid
In principle, it should be possible to find the density of a volume of strain confined electron-hole liquid quite accurately with Alfvén wave resonance experiments and simultaneous imaging experiments. As was shown in Chapter I, the Alfvén wave resonant field, $H_c$ has approximately the following relationship with the radius $R$ and liquid density $n$ for a spherical drop

$$H_c^2 \propto nR^2.$$  \hspace{1cm} (III-1)

I therefore set up the experimental apparatus shown schematically in Fig. III-1 for simultaneous imaging and Alfvén wave resonant experiments. The Alfvén wave absorption apparatus was set up to be the same as for the buildup experiments discussed in Chapter II. The pump light source was a mechanically chopped 2W CW argon ion laser. The laser beam was directed through a hole in a front surface mirror, through a camera lens, and onto the bottom of the sample. The luminescence was then collected from the bottom of the sample through the same lens, imaged onto two crossed movable slits and collected with a cooled Ge photodiode. The luminescence is separated from the laser light by the front surface mirror with a hole in it and by a Corning No. 7-56 infrared transmitting filter. Due to poor luminescence collection efficiency, it was necessary to remove the spectrometer in order to perform infrared imaging. However, it was found that the only source of infrared radiation was from the SCEHL by checking the luminescence with an interference filter centered at 1.748 $\mu$m with .01 $\mu$m width. It was not practical to perform imaging with this filter however since it had to be rotated to be sensitive at the SCEHL wavelength, $\sim$ 1.753 $\mu$m, and had 40% peak transmission.

Typical imaging results are shown in Fig. III-2. We observed the surprising result that the EHL deforms markedly in a magnetic field.
We had discovered the magnetostriction of the EHL III.2-III.4. This effect was independently discovered by Störmer and Bimberg III.5-III.7 while searching for a new phase transition of the liquid III.8.

In an applied magnetic field the energies of the four conduction band valleys can shift with respect to one another because the cyclotron masses will, in general, be different for each valley. In principle, as suggested in references III.5 - III.7, this could modify the shape of the potential well by changing the relative population of the valleys. In the experimental condition $\hat{H} \perp \langle 111 \rangle$, this shape change could occur only if more than one conduction band is occupied by electrons within the SCEHL. This is because, for $\hat{H} \perp \langle 111 \rangle$, a magnetic field further raises the three unoccupied valleys in energy with respect to the occupied valley. EHD are attracted by an octahedral strain maximum only if the strain is large enough that just one conduction band valley is occupied as was shown in Chapter I and has been experimentally confirmed as discussed in Chapter V. For example, in Fig. III-3b the shaded regions are the only parts of the crystal where more than one valley is occupied. Hence, until the SCEHL is large enough to fill the unshaded region in zero field, a magnetic field applied perpendicular to the stress direction would have no effect on the shape of the drop, contrary to the observation in Fig. III-2. This is further illustrated theoretically in Fig. III-4 which shows EHD energy contours recalculated as in Fig. III-3b, but now with a 50 kOe magnetic field applied parallel to the $\langle 112 \rangle$ direction. As expected, the shape of the constant energy contours is affected only in the region of the crystal shaded in Fig. III-3b. Thus this static bandshifting mechanism cannot explain the experimental results where a large distortion is observed even for $H < 5$ kOe. It should be noted however, that if there are two spatially separate
drops in the crystal, the electrons will occupy a different \( \text{(111)} \) valley in each as discussed in Ref. III.14. Consequently, a magnetic field can alter the relative populations of carriers in the two drops by this mechanism as observed in reference III.6.

Markiewicz III.3, III.9, III.10 has developed a detailed dynamic model for this effect. The source of the energy causing the shape change is the recombination current

\[
\mathbf{J} = \mathbf{n} \mathbf{v}, \quad (\text{III-2})
\]

where \( \mathbf{v} \) is the center of mass velocity field, which is necessary for an EHD to exist in steady state. The current is governed by the continuity equation

\[
\mathbf{v} \cdot \mathbf{J} = -\frac{n}{\tau_0} \quad (\text{III-3})
\]

where \( \tau_0 \) is the pair recombination time as was first shown by Kaminskii and Pokrovskii. III.11 The current is supplied by photoproduced droplets, excitons, and carriers which are created at the crystal surface and flow into the strain well as discussed in Chapter II. For a spherically symmetric flow into the surface, the recombination currents in zero field are shown in Fig. III-5a for this case the drift velocity is

\[
\mathbf{v} = -\frac{\mathbf{r}}{3\tau_0}, \quad \text{which for a carrier near the surface of a drop with radius } R = 200 \mu m \text{ and lifetime } \tau_0 = 500 \mu s \text{ is approximately } 15 \text{ cm/s.}
\]

In a magnetic field the recombination currents are deflected by the Lorentz force, as shown in Fig. III-5b. Because electrons and holes are deflected in opposite directions, there will be a net azimuthal electrical current about the field direction. This current gives rise to a net paramagnetic (i.e., positive) magnetization \( M \), the "recombination magnetism." III.11 For a spherical droplet, the
contribution to the magnetization by either electrons or holes is given by

\[ M_R = 4\pi n_o e\omega_c^2 R^3/45c\tau_o \]  

(III.4)

where \( \tau \) is the carrier scattering time, and \( \omega_c = eH/m_c c \) is the cyclotron frequency, and \( m_c \) is the cyclotron effective mass.

The distribution of recombination currents and hence the magnetization depends upon the shape of the drop. The total energy of the drop can be lowered if the drop distorts in such a manner as to increase \( M_R \), since the magnetic energy is

\[ E_m = -\frac{1}{2} M \cdot \vec{H}_o \]  

(III-5)

where \( \vec{H}_o \) is the external applied field, and \( \vec{M} \) is the total magnetization. The potential energy of e-h pairs in the strain well \( E_s(r) \) acts as a restoring force; thus, the minimization of the total energy \( E_m + E_s \) determines the drop shape in a magnetic field. In general, the recombination magnetism is increased by increasing the radial currents perpendicular to the field. This corresponds to a flattening of the drop along the field direction, as is generally observed experimentally. Other contributions to the magnetization of the drop such as the usual Landau diamagnetism are largely volume effects and depend only very weakly on the shape. Thus even if the total magnetization is diamagnetic, the drop will change shape as if the recombination magnetization were the only effect.

As a consequence of this dynamic mechanism, we expect the drop flattening can be varied by moving the focused laser spot on the crystal surface, because the recombination current flow pattern depends on the carrier flux distribution at the drop surface. For example, if all the carriers enter the drop through a small area of
surface lying on the drop diameter parallel to the magnetic field (Fig. III-5c and d), then the recombination current perpendicular to \( \mathbf{H}_0 \) will be predominantly outward, from the center toward the surface. In this case, the resultant recombination magnetism is diamagnetic, and the drop should elongate along the magnetic field, reducing the recombination currents perpendicular to the field. As was seen in Chapter II, this channeling along the field can be accomplished by placing \( \mathbf{H}_0 \parallel \langle 111 \rangle \parallel \) stress and pumping below the drop location.

These ideas explain the low field behavior of an EHD in a magnetic field and give a qualitative idea as to the effects involved. In order to explain the saturation of the effect at high field and give detailed predictions of the drop shape in decay, Markiewicz developed a more sophisticated magneto-hydrodynamic model for the effect. \(^{\text{III.9,III.10}}\) The basic equations of the carrier flow for a cylindrical coordinate system with directions \( \hat{z} \parallel \mathbf{H}_0 \), \( \hat{\rho} \), and \( \hat{\phi} \) are Eq. (III-2) and

\[
\begin{align*}
\nabla \mathbf{p} &= -2\sigma_0 \frac{H_0^2}{c^2} \mathbf{v}_\rho \mathbf{\hat{\rho}}. 
\end{align*}
\]

The strain well is taken to be parabolic: the strain energy per pair is \( \alpha r^2 \) where \( r \) is the radial distance from the center of well. The other symbols are defined as follows: \( \mathbf{v}_\rho \) is the \( \rho \) component of the velocity; \( \mathbf{p} \) is the pressure, and \( \sigma_0 \) is the zero field conductivity of the EHL. Viscous effects are negligible. To these equations is added an equation of state for the liquid which in general can be calculated from the thermodynamic relation \( p = n^2 \frac{\partial E(n,T)}{\partial n} \) and calculated curves of pair energy, \( E \), versus density and temperature (see, for instance, Ref. III.11). However, Markiewicz found that the much simpler case of incompressible liquid could be used at all but the very highest fields. Therefore
where \( n_0 \) is the equilibrium density.

There are also two boundary conditions. First, the pressure must balance the surface tension, but, because the surface tension is negligible for the SCEHL, this becomes

\[
p = 0 \quad \text{(at surface)}.
\]  

(III-8)

The second boundary condition states that the normalized component of the flux of particles into the drop is a known function depending on the flux of carriers outside the drop characterized by \( \vec{w} \):

\[
\hat{n} \cdot (n\vec{v}) = \frac{n_0}{\tau_0} \hat{n} \cdot \vec{w} \quad \text{(at the surface)}
\]  

(III-9)

where \( \hat{n} \) is a unit vector perpendicular to the surface.

Markiewicz then showed that although the pumping in general could be quite anisotropic, any anisotropies in pumping perpendicular to the field could be replaced by their average value for the calculation. Therefore he used symmetric pumping for simplicity

\[
\vec{w} = -R_0^2 \hat{n} \times 3r^3
\]  

(III-10)

where \( R_0 \) is a measure of the pumping intensity; the drop will be spherical with radius \( R_0 \) in zero field. These equations can be combined into equations for the surface of the drop and solved numerically. The results using the values for the parameters tabulated in reference III.9 show an appreciable flattening at 2 kOe and an oblateness saturation of \( R_\perp/R_\parallel \approx 2 \) where \( R_\parallel \) is the drop dimension parallel to \( \vec{H} \) and \( R_\perp \) is the drop dimension perpendicular to \( \vec{H} \), at about 10 kOe in steady state independent of initial drop size for the SCEHL. Further he found that during the decay in high magnetic fields \( R_\perp \) is approximately constant. The oblateness therefore increases in time;

\[
R_\perp/R_\parallel \propto \exp \left( t/\tau_0 \right).
\]  

(III-11)
In order to more thoroughly study the effect of magnetic field on SCEHL, two types of infrared imaging experiments were performed. First, in order to get a good qualitative understanding of the effect, photographic experiments using an infrared vidicon as shown in Fig. III-6 and discussed in detail in references III.12 - III.14 were performed. The sample was pumped with a 2W CW argon ion laser which was focused onto the sample with a simple lens mounted on a micrometer driven x-y-z translation stage. This apparatus could be used to reproducibly move the laser spot on the sample. The luminescence from the SCEHL was focused on the imaging surface of an infrared vidicon, displayed on a television monitor and photographed.

Typical results from these photographs are shown in Figs. III-7 through III-9. Figure III-7 shows that the effect is largely symmetric about the field direction except where the strain well departs significantly from parabolic near the plunger. It can also be seen that a flattening of about 2:1 or 3:1 parallel to the applied field is observed.

Figure III-8 shows that contrary to the predictions of the static band shifting modelIII.5-III.7 a distortion, in fact a distortion of similar magnitude, is observed when $\mathbf{H}||\langle 001\rangle$. As explained in reference III.5, a field parallel to $\langle 001\rangle$ shifts all electron valleys equally in energy, and therefore no shape change should occur when $\mathbf{H}||\langle 001\rangle$ in the static band shifting model.

Finally, Fig. III-9 shows that the expected change in shape with changing pump position for $\mathbf{H}||\langle 111\rangle$ stress is observed. In these successive photographs the laser spot is translated from approximately on top of the drop in Fig. III-9a to about 1mm below the drop location for III-9d. In all cases the drop is symmetric about the strain,
and magnetic field axis as confirmed by photographs of the bottom view. The drop has gone from oblate shape (paramagnetic $M_R$) to prolate shape (diamagnetic $M_R$) just as predicted by Markiewicz's theory with the data of Chapter II, contrary to any static model. III.4, III.9, III.10

Next, more quantitative results were obtained by performing the image scanning experiments shown in Fig. III-10. The sample was pumped as for the photographic experiments except the laser was chopped with the special high speed chopper described in reference III.12. The luminescence was collected with a camera lens and imaged onto the front slit of the spectrometer. A special crossed slit could be inserted into the spectrometer to allow spatial resolution in two dimensions. The image was scanned across the spectrometer with the special stepper motor-driven two dimensional scanning apparatus which I designed. A block diagram of this apparatus is shown in Fig. III-11.

Figure III-12 shows typical data of drop dimensions as a function of magnetic field intensity. These data were obtained by translating the image of the SCEHL across a narrow slit with the special scanning apparatus and measuring the transmitted luminescence intensity. In this way a spatial profile of the drop luminescence parallel and perpendicular to the magnetic field was obtained, giving the radial extent of the droplet in these two directions. The data here were taken with relatively high light level, absorbed power $P_a \approx 90$ mW, and moderate applied force 9 kgf. This force is applied across a contact area $\approx 1 \text{mm}^2$ as in reference III.14. The principal results are that a) the primary change in shape occurs for $H \lesssim 10$ kOe, and that b) above $H \approx 10$ kOe the amount of distortion appears to saturate with increasing field. This saturation field depends on the magnitude
of the applied stress and weakly on the drop size (see Figs. III-13 and III-15). This behavior of shape with magnetic field is just as expected from Markiewicz's theory. III.4, III.9, III.10

The effect of changing the applied stress is displayed in Fig. III-13. In a given magnetic field the distortion for 4.5 kgf is considerably greater than for 18 kgf. Not surprisingly, the effect of magnetic field on drop shape is less for a deep potential well than for a shallow one because the strain gradient acts as a restoring force. For the samples reported on here the luminescence decay time is measured to be \( \approx 500 \mu s \). Figure III-14 shows that about 20% variation in this lifetime is observed for applied fields up to 25 kOe.

Bimberg suggests that, even though only one conduction band minimum is occupied at the center of the strain well, it is possible that near the surface of the drop the strain is sufficiently reduced that several minima are occupied. While the calculations of Fig. III-3b and III-4 indicate that this is unlikely, it is possible to give further direct experimental evidence. If multiple valleys are occupied only near the surface of the SCEHL, then the flattening effect should be much less pronounced for a smaller drop, and should vanish for drops of radius \( R < R_{\text{min}} \), for which all electrons are in single conduction bands at \( H = 0 \). Assuming for the moment that the distortion were due to static bandshifting, a reasonable estimate for \( R_{\text{min}} \) should be the saturated value of \( R \) at high fields, which for the data of Fig. III-12 would yield \( R_{\text{min}} \approx 300 \mu m \). Figure III-15 shows that this is not the case. For the same sample and stress condition as in Fig. III-12 the laser pump intensity was reduced until the initial drop radius was \( \approx 150 \mu m \). Clearly, the oblateness has not changed significantly. In fact, it is
about the same as is expected from the results of Markiewicz's
theory.

Finally, the time dependence of the shape was examined. Because
carriers move much more easily parallel to a magnetic field than per-
pendicular to it, the decay of the SCEHL is greatly influenced by a
magnetic field. Figure III-16 shows the variation of the drop radius
as the SCEHL decays. The data are obtained by a time resolved imaging
technique. We use basically the same slit-scanning

technique as in Figs. III-12, III-13, and III-15; however, by means
of gated boxcar integration the luminescence profile is measured
at discrete times after the excitation laser is chopped off. As
was seen in Chapter II, within a few microseconds after the laser is
chopped off, nearly all of the non-equilibrium carriers are contained
within the drop, and no net particle flow into the drop remains.
Hence the spatial decay of the isolated strain-confined liquid in a
magnetic field is determined. The experiment was performed with the
stress and \( \mathbf{H} \) orientation of Fig. III-9, for the excitation conditions
of both Fig. III-9a and III-9d. The data in Fig. III-16 show that the
dimension of the drop parallel to the magnetic field decays with \( \tau_0 \)
consistent with the data of Fig. III-14 and the predictions of Markiewicz,
Eq. (III.11). This is true independent of the initial drop shape before
the light is removed. It is clear that the shape of the electron-hole
drop depends upon the dynamic particle currents within the drop which
are altered by the magnetic field.

Thus the dynamic magneto-hydrodynamic model of the shape of the
SCEHL in a magnetic field of Markiewicz has been confirmed in detail.
There is remarkable agreement between experiment and theory in all
cases.
CHAPTER III: References


CHAPTER III: Figure Captions

Fig. III-1 Apparatus for simultaneous imaging and Alfvén wave resonance experiments.

Fig. III-2 Typical results obtained from the imaging experiment in a magnetic field. Drop radius versus magnetic field is plotted parallel and perpendicular to the field $R_y \parallel \vec{H}; R_z \perp \vec{H}$. Note that there is some overall loss of efficiency and that the drop flattens parallel to the applied field. The insert shows the experimental configuration. The sample is permanently stressed; $T = 1.8$ K. The solid lines are smooth curves through the data. Sample CR14.

Fig. III-3 a) Photograph of the SCEHL in a $4 \times 4 \times 2$ mm$^3$ crystal of ultrapure dislocation-free Ge. b) Calculated lines of constant electron-hole liquid energy (meV) in a square Ge sample, $\langle 110 \rangle$ face, which is stressed at the top by a nylon plunger with 9 kgf. The plunger has a 6 mm radius of curvature. The applied stress direction is $\langle 111 \rangle$. This strain pattern is obtained from a finite element computer analysis by Markiewicz. The EHL is confined to a point of maximum shear strain.

Fig. III-4 Numerical calculation of the EHL energy shift (meV) for the geometry of Fig. III-3b and an applied magnetic field $H = 50$ kOe: $\vec{H} \parallel \langle 112 \rangle$. Note that there is no change in the shape of the contours outside the shaded region of Fig. III-3b.

Fig. III-5 a) Schematic diagram showing recombination currents of electrons (solid arrows) and holes (dashed arrows) in zero magnetic field. b) Modification of these drift currents caused by the Lorentz forces in an applied field, assumed
parallel to the z-axis. A net circulating current results, leading to a paramagnetic "recombination magnetization" and the distortion in drop shape usually observed.

c) Schematic of a non-radial distribution of recombination currents which would be expected from the results of Chapter II with \( \hat{\mathbf{z}} \parallel (111) \), similar to c) but in a non-zero field. In this case, the induced electric current produces a diamagnetic magnetization.

Fig. III-6 Apparatus for recording an image of the spatial distribution of the recombination luminescence from the SCEHL in Ge.

Fig. III-7 a) Drawing of the crystal mounting and orientation. Both face and edge views are shown. b) Image from the infrared vidicon showing a single EHD within the crystal. Laser power absorbed \( P_a \approx 50 \text{ mW} \); \( T = 1.7 \text{ K} \). The contact force is \( \approx 9 \text{ kgf} \); and \( H = 0 \). c) Distortion in drop shape when \( H = 20 \text{ kOe} \parallel (11\bar{2}) \) axis. All other conditions as in b).

d) Distortion in drop shape when \( H \parallel (1\bar{1}0) \). All other conditions as in c). Sample CR37.

Fig. III-8 a) Diagram of stress and field directions for \( H \parallel (001) \) experiment. b) Image of the drop luminescence. \( H = 0 \), and \( P_a \approx 50 \text{ mW} \). b) Change in shape observed when a 20 kOe field is applied along the (001) axis. The sample is permanently stressed. III.14 This demonstrates irrefutably that the static band shifting theory III.5-III.7 is incorrect. Sample CR36.

Fig. III-9 Change in EHD shape observed as the laser spot is translated. \( P_a \approx 50 \text{ mW} \); \( H \parallel (111) \); \( H = 20 \text{ kOe} \); \( T = 1.7 \text{ K} \). Laser position is a) 0.5, b) 0.9, c) 1.1, and d) 1.5 mm from the top
of the crystal which is stressed, defined by the bright horizontal line. All of these drops were observed to be cylindrically symmetric about the field direction. This is a dramatic confirmation of Markiewicz's magneto-hydrodynamic theory.

Fig. III-10 Schematic view of the experimental apparatus used in imaging experiments. For steady state experiments the boxcar is replaced by a lock-in.

Fig. III-11 Detailed schematic of scanner control and deflection mirror.

Fig. III-12 Field dependence of the drop distortion for 9 kgf applied. \( R_\parallel \) is the EHL dimension along \( H \) and \( R_\perp \) its dimension perpendicular to \( H \). This dimension is extracted from a luminescence image scan by taking the full width at half maximum of the scan and dividing by \( \sqrt{2} \) as would be expected for a spheroid. For details see Ref. III.4 and III.13. These data correspond to the geometry of Fig. III-7d. \( P_a \approx 90 \text{ mW} \). Solid lines are smooth curves through data.

Fig. III-13 Field dependence of the drop distortion for two other values of the applied force: triangles, 4.5 kgf, and circles, 18 kgf. As in Fig. III-12, \( H \parallel \langle 110 \rangle \). \( P_a \approx 3 \text{ mW} \). Solid lines are smooth curves through data.

Fig. III-14 Weak dependence of the SCEHL luminescence decay time as a function of \( H \). \( P_a \approx 3 \text{ mW} \); the applied force is 9 kgf. The solid line is a smooth curve through the data.

Fig. III-15 Field dependence of the drop distortion for the same experimental conditions as in Fig. III-12 except \( P_a \approx 3 \text{ mW} \). The drop distortion is present, even though the initial radius is smaller than the saturation value of \( \approx 300 \mu\text{m} \) for \( R_\parallel \) in Fig. III-12.
Fig. III-16 Time resolved decays of the EHD in a magnetic field $H = 20 \text{ kOe}$, $H \parallel (111)$, for two different initial drop shapes. The light ($P_a \approx 90 \text{ mW}$) is chopped off at $t = 0$, and a time resolved slit scan gives the EHD dimensions parallel ($R_\parallel$) and perpendicular ($R_\perp$) to the field. It can be seen that $R_\perp$ stays constant, while $R_\parallel$ decays with some characteristic time constant $\tau_0$. a) Laser excitation as in Fig. III-9a. $\tau_0 = 560 \mu s$. b) Laser excitation as in Fig. III-9d. $\tau_0 = 660 \mu s$. 
Fig. III-1
Fig. III-2
Fig. III-3
EHL ENERGY SHIFT
(meV)

<111> STRESS
(110) FACE
H=50 kOe

Fig. III-4
Fig. III-6
Fig. 11-7

(a) stress
(b) H=0
(c) H → H0
(d) H=20kOe

Ge crystal

hole

mirror

45°
Fig. III-11
Fig. III-16
CHAPTER IV

The Phase Diagram of the Strain-Confined Electron-Hole Liquid
The strain-confined electron-hole liquid provides a good medium for the study of the equilibrium properties of the EHL in the lower degeneracy configurations of Ge, Ge[1:2] and Ge[1:1], because the intrinsically long lifetime of \( \sim 500 \mu s \) is not reduced due to surface recombination or boil-off. The surface recombination is nonexistent for the SCEHL because the liquid is in the interior of the crystal, and the boil-off is inhibited because all carriers and FE emitted from the EHL are attracted back into the liquid by the strain gradient in a time short compared to their lifetime. However, the strain well must be taken into account in any of these studies.

Assuming thermal and diffusive equilibrium, the chemical potential \( \mu \) of the system must be constant within the strain potential well. The strain energy is approximately parabolic in the distance \( r \) from the center of the well:

\[
E_s = \alpha r^2 \tag{IV-1}
\]

where the well parameter

\[
\alpha \approx 10 \text{ meV/mm}^2 \tag{IV-2}
\]

in our experiments. Thus in the potential well

\[
\mu = \mu_0 + \alpha r^2 \tag{IV-3}
\]

where \( \mu_0 \) is the chemical potential of the species in question in the absence of the strain well.

Assuming the FE are an ideal gas, the unperturbed exciton chemical potential is

\[
\mu_{ox} = k_b T \ln (n_x V_Q) \tag{IV-4}
\]

where \( n_x \) is the density of excitons, \( k_b \) is Boltzmann's constant and
is the exciton quantum volume. The degeneracy and density of states mass of the exciton are \( v_x \) and \( m_x \), respectively. In the strain well the density \( n_x \) becomes position dependent. As was shown in Chapter I, above a small threshold stress \( \sim 3 \text{ kgf/mm}^2 \) the FE binding energy is stress independent. It is therefore straightforward to get the FE density as a function of position from the above equations:

\[
  n_x(r) = n_{xo} \exp \left( -\frac{ar^2}{k_bT} \right) \quad (IV-6)
\]

where \( n_{xo} \) is the FE density at the bottom of the well. Figure IV-1 from Ref. IV.3 shows that the distribution of excitons in the strain well actually has this form. This data also confirms that FE and EHL are both attracted into the strain well and that the SCEHL forms at the bottom of the well. From Eq. (IV-6) it can be seen that the halfwidth of the FE distribution in the well should be proportional to \( T^{3/2} \). This behavior has been verified by Gourley and Wolfe in Si.

When the FE luminescence is collected from the entire strain well, then the luminescence intensity

\[
  I_x \propto N_x \quad (IV-7)
\]

Assuming \( ar^2 \gg k_bT \) the total number of FE in the well is

\[
  N_x = 4\pi \int_0^\infty r^2 n_x(r) \, dr = \frac{1}{2} (k_bT/a)^{3/2}, \quad (IV-8)
\]

and therefore

\[
  I_x \propto T^{3/2} n_{xo} \quad (IV-9)
\]

This last relation is needed later to interpret experimental results.

Within the SCEHL a similar argument about effects of the strain well is used. At \( T = 0 \) the unperturbed EHL chemical potential is
\[ \mu_{\lambda} = E + \frac{p}{n_{\lambda}} \]  

(IV-10)

where \( E \) is the energy per pair, \( p \) is the fluid pressure, and \( n_{\lambda} \) the liquid density. For small deviations from the equilibrium density \( n_0 \) and for zero of energy at the FE energy,

\[ E = -\phi + \frac{1}{2} E''(n_{\lambda} - n_0)^2, \]  

(IV-11)

where \( \phi \) is the EHL equilibrium pair condensation energy, and \( E'' \) is related to the compressibility: IV.5, IV.6

\[ E'' = -\left. \frac{\partial^2 E(n_{\lambda})}{\partial n^2} \right|_{n_{\lambda} = n_0}. \]  

(IV-11a)

The pressure is then given as

\[ p = n^2 \frac{\partial E}{\partial n} = n^2 E''(n_{\lambda} - n_0). \]  

(IV-12)

Keeping only terms linear in \( n_{\lambda} - n_0 \), Eq. (IV-3) becomes

\[ \mu_{\lambda} = -\phi + n_0 E''(n - n_0) + \alpha r^2. \]  

(IV-13)

At the surface of the drop \( p = 0 \) assuming that the strain energy is much larger than the surface energy. Thus, for a drop of radius \( R \) in the strain well

\[ n(r=R) = n_0, \]  

(IV-14)

and

\[ n(r) = n_0 \left[ 1 + \hat{\alpha}(R^2 - r^2) \right] \]  

(IV-15)

where \( \hat{\alpha} = \alpha/(n_0^2 E'') \). Markiewicz and Kelso IV.5 find \( \hat{\alpha} \approx 16 \, \text{mm}^2 \) for \( \alpha = 10 \, \text{meV/mm}^2 \). Therefore for a small drop, \( R < 80 \, \mu\text{m} \), the EHL density will be within 10% of \( n_0 \) throughout the drop at \( T = 0 \).

Finally including both EHL and FE in the strain well and setting \( \mu_x = \mu_{\lambda} \) and noting that the strain energy is the same for both species one obtains

\[ k_b T \ln \left( n_x Q_x \right) = -\phi + n_0 E''(n_{\lambda} - n_0). \]  

(IV-16)

Also setting the pressures equal at \( r = R \) and noting from the equation of state for an ideal gas \( p_x = n_x k_b T \), I obtain from Eq. (IV-12)
to first order in $n_x - n_o$

$$n_x k_b T = n_o^2 E_0 (n_x - n_o).$$  \hspace{1cm} (IV-17)

Eliminating $n_x$ in Eq. (IV-16) results in

$$- \frac{\phi}{k_b T} = \ln \left( \frac{n_x V Q_x}{n_o} \right) - \frac{n_x}{n_o}.$$  \hspace{1cm} (IV-18)

At low temperatures, $k_b T \ll \phi$ and $n_x / n_o \approx 0$. Along with the results of Eq. (IV-9), this equation allows a thermodynamic determination of the SCEHL condensation energy $\phi$.

Next I will consider the effect of the strain potential well on the characteristic luminescence spectral lineshapes of the EHL and FE. The EHL lineshape is simply a convolution of the density of occupied electron and hole states. \hspace{1cm} (IV.8)

This leads to the EHL luminescence lineshape \hspace{1cm} (IV.9,IV.10)

$$I_x (h\nu) = \int D_e f_e f_h \delta (E_g + E_e + E_h - h\omega_{ph} - h\nu) dE_e dE_h,$$  \hspace{1cm} (IV-19)

where $D_e$ and $D_h$ are the electron and hole densities of states, $f_e$ and $f_h$ are the respective Fermi distribution functions, $E_g$ is the energy of the indirect gap, $h\omega_{ph}$ is the energy of the phonon emitted along with the photon and the integral is taken over all electron and hole energies. This makes the full width of the line equal to the sum of the Fermi energies of the electrons and holes which is the Fermi energy of the liquid

$$E_F = E_F^e + E_F^h.$$  \hspace{1cm} (IV-20)

As it was shown in Eq. (IV-15), for small volumes of SCEHL the $n_x$ and therefore $E_F$ are approximately constant in position. This has been experimentally verified by Wolfe et al. \hspace{1cm} (IV.3) as shown in Fig. IV-2. It can be seen that for low powers the linewidth is constant as expected. Therefore it is expected that the lineshape, Eq. (IV-19), will give
good results for the Fermi level $E_{F,L}$ and density of small volumes of SCEHL. The Fermi level of the liquid as observed in luminescence is defined as

$$E_{F,L} = E_{go} - E_x - \phi - \hbar \omega_{ph}$$  \hspace{1cm} (IV-21)

where $E_{go}$ is the energy of the gap at the bottom of the strain potential well;

$$E_{go} = E_g (r=0).$$

The shape of the FE line can also be easily calculated. Elliot\textsuperscript{IV.11} showed that the absorption edge for FE in Ge at low temperatures is proportional to their density of states $D_x(E)$. Therefore the luminescence lineshape for FE with binding energy $E_x$ is

$$I_x(h\nu) = D_x(E) \exp \left[-(\hbar \nu - E_x + \hbar \omega_{ph})/T\right].$$  \hspace{1cm} (IV-22)

As was first pointed out by Wolfe\textsuperscript{IV.12} $D_x(E)$ must be recalculated for excitons in the strain well. The FE may be considered free particles in a three-dimensional harmonic oscillator potential $\alpha r^2$. From simple quantum mechanics this leads to a series of energy levels; the $n$th level has energy

$$E_n = (n + \frac{1}{2})\hbar \omega_s$$  \hspace{1cm} (IV-23)

and degeneracy

$$\gamma_n = n(n + 1)/2$$  \hspace{1cm} (IV-24)

where

$$\omega_s = (\alpha/2m_x)^{1/2}$$  \hspace{1cm} (IV-25)

and $m_x$ is the exciton density of states mass. For the strain potential well in Ge $\alpha \approx 10$ meV/mm$^2$ and $m_x = 0.4m_0$;\textsuperscript{IV.14} this gives $\hbar \omega_s \approx 50\mu K << T$. Therefore the continuum limit is appropriate and

$$D_x(E) \propto (E_0 - h\nu)^2$$

and from Eq. (IV-22)

$$I_x(h\nu) \propto (E_0 - h\nu)^2 \exp \left[-(E_0 - h\nu)/T\right]$$

where
\[ E_0 = E_{go} - E_x - \hbar \omega_{ph}. \]

This lineshape has been confirmed in Si by Wolfe at al. IV.12

The apparatus for the experiments discussed here is shown schematically in Fig. IV-3. The 4x4x2 mm ultrapure Ge sample was stressed with a Hertzian contact stress as discussed in Chapter I and placed in an optical liquid He cryostat. The pumping light was a chopped 2 W CW Ar ion laser which was attenuated with Corning type 7-98 neutral density filters to obtain the low light powers necessary. The luminescence was collected, spectrally resolved with a Jarrel-Ash ¼ m spectrometer, and detected with a cooled Ge photodiode. IV.13 The output of the detector was amplified and then analyzed with a PAR model 128A lock-in amplifier. The output of the lock-in was recorded on a strip chart recorder.

Typically spectra for different laser powers absorbed \( P_a \) are shown in Fig. IV-4. At the lowest powers (Fig. IV-4a) only FE are observed, and their luminescence lineshape agrees quite well with the expected shape from Eq. (IV-25) convolved with the spectrometer slit function (open circles in Fig. IV-4a). At slightly higher powers (Fig. IV-4b) both EHL and FE are observed. As the power is increased above this EHL threshold, the EHL luminescence intensity grows quickly, while the FE luminescence intensity is approximately constant. Figure IV-4c shows this. Here the EHL lineshape is fit with Eq. (IV-9) including the necessary complications due to finite temperature Fermi functions and stress dependent hole masses. IV.3, IV.6 This fit is shown by the dots in Fig. IV-4c and corresponds to the following density and Fermi energy at \( T = 3.3 \text{ K} \): \( n_0 = (4.6 \pm 0.3) \times 10^{16} \text{ cm}^{-3} \); \( E_F = 4.3 \pm 0.3 \text{ meV} \); and \( E_F^0 = 2.1 \pm 0.10 \text{ meV} \). The difference between \( E_0 \) and the liquid Fermi level
\[ E_{\text{EHL}} \] gives a spectroscopic determination of \( \phi \), the condensation energy of the EHL, which is \( \phi = 10.8 \pm 1 \text{ K at } T = 3.3 \text{ K} \). When this procedure for getting \( n_o(T) \) and \( \phi(T) \) was attempted at higher \( T \), I found that the apparent \( n_o(T) \) was increasing, and the apparent \( \phi(T) \) was getting smaller as \( T \) increased. This is in direct contradiction with the known behavior for these quantities especially \( n_{\text{IV.8,IV.13,IV.15-IV.17}} \)

\[ n_o(T) = n_o(0)(1 - n_o'T^2) \]  

(IV-27)

where \( n_o' \) is the first correction to \( n_o \) from Fermi liquid theory and is positive. This discrepancy is explained by noting compressibility of the liquid is getting larger as the temperature increases. Therefore the assumption that \( T = 0 \) used in getting Eq. (IV-15) and the estimates of the importance of strain thereafter, are no longer valid and a determination of \( n_o(T) \) for these higher temperatures is not possible from these data. The determination of \( \phi(T) \) is invalidated because \( E_{\text{EHL}}(T) \) is not properly found for the compressed SCEHL and thus using this spectroscopic measurement the apparent \( \phi \) is lower than the actual \( \phi \) at higher \( T \).

An independent measure of the SCEHL condensation energy can be made by measuring thresholds for EHL formation in the strain well. Typical data are shown in Fig. IV-5. At various laser pump levels the total luminescence intensity is recorded for the FE and EHL. The EHL intensity is then linearly extrapolated to \( I_x = 0 \), and \( I_x \) at that power is recorded at various temperatures. This is the free exciton luminescence intensity \( I_{\text{XO}} \) at the SCEHL threshold. Because the FE density is highest at the bottom of the well, this is where the SCEHL will nucleate (see Fig. IV-1). Therefore using Eqs. (IV-9) and (IV-18), \( I_{\text{XO}}(T) \) can be related to \( \phi \) giving
\[- \frac{\phi}{T} = \ln \left( \frac{I_{xo}}{T^3} \right) + \text{const.} \tag{IV-28} \]

The data is plotted in this way, $I_{xo}/T^3$ versus $1/T$, on semilog paper and displayed in Fig. IV-6. The solid line shows a good fit to the data for $\phi = 9.2 \pm 1$ K which is in good agreement with the spectroscopic determination.
CHAPTER IV: References


CHAPTER IV: Figure Captions

Fig. IV-1 Spatial profiles of luminescence intensity taken with a mirror to show the distribution of a particular phase through the sample as shown in Fig. I-7c. The laser is incident on the face at $x = 0$. Spatial resolution $\approx 80 \mu m$; $T = 4.2 K$. (a) FE gas phase luminescence near threshold for EHL formation. The open circles are a fit using Eq. (IV-6) confirming the expected spatial distribution with $a = 11 \text{meV/mm}^2$. (b) Liquid phase luminescence intensity at somewhat higher excitation level. Notice that it is in the same part of the crystal as the FE gas as expected. Sample CR38.

Fig. IV-2 Comparison of the full width at half maximum linewidth $\Delta E$ for the LA-assisted luminescence from the EHL in unstressed Ge and the SCEHL. At low powers both $\Delta E$ are constant as expected for a constant-density liquid phase. At $P_a \approx 5 \text{mW}$, corresponding to $R \approx 200 \mu m$, the SCEHL linewidth becomes noticeably broadened. $T = 1.8 K$. Sample CR38.

Fig. IV-3 Experimental apparatus for threshold studies.

Fig. IV-4 Luminescence spectra for strain confined phases; $T = 3.3 K$ (a) Absorbed laser power $P_a = 0.05 \text{mW}$. Only the FE luminescence is seen. Open circles are a fit to the line using Eq. (IV-25); the only adjustable parameters are the height and position of the line. (b) $P_a = 0.14 \text{mW}$. Both FE and SCEHL are observed. (c) $P_a = 0.42 \text{mW}$. The SCEHL luminescence has increased markedly while the FE luminescence is about the same. The dots are a fit to the SCEHL lineshape using the procedure outlined in the text. The
procedure for the spectroscopic determination of the condensation energy is shown schematically; \( E_0 - E_{FL} = \phi \).

Sample CR36.

**Fig. IV-5** Luminescence intensity versus laser power for the FE and EHL in the strain well. The threshold can be easily seen in both curves. The solid curves are straight lines through the data. Sample CR36.

**Fig. IV-6** Semilog plot of \( \frac{I_{X0}}{T^3} \) versus \( 1/T \) used to determine \( \phi \). Solid line is the best fit for \( \phi = 9.2 \) K. Sample CR36.
Fig. IV-1
Fig. IV-2
LASER POSITIONING MIRRORS

7-98 FILTERS

SAMPLE

STRESSING ROD

7-56 FILTER

DEWAR WINDOWS

COLLECTING LENS

FOCUSING LENS

TO REF. ON LOCK-IN

Fig. IV-3
\[ T = 33K \]

Stress = 9 Kg along \( \langle \text{III} \rangle \)

Fig. IV-5
$I_{x_0}/T^3$ (Rel. units)

$\phi = 9.2 \text{ K}$

Fig. IV-6
CHAPTER V

Magneto-Oscillations in the Luminescence from the Strain Confined Electron-Hole Liquid
Studies of the luminescence from the electron-hole liquid in a magnetic field have proven to be most fruitful.\textsuperscript{5.1} Magneto-oscillations similar to the de Haas-van Alphen or Shubnikov-de Haas effects periodic in $1/H$ were first observed in the luminescence of the EHL in unstrained Ge by Bagaev et al.\textsuperscript{5.2} and explained as being due to a change in the binding energy with magnetic field as the Landau levels pass through the Fermi level. Keldysh and Silin\textsuperscript{5.3} performed a calculation which showed that these oscillations were due to density variations rather than binding energy variations. This was shown conclusively experimentally by Betzlar et al.\textsuperscript{5.4-5.9} and Martin et al.\textsuperscript{5.10} These experiments have lead to the conclusive demonstration of the metallic nature of the EHL,\textsuperscript{5.1} the most consistent measurement of the quantum efficiency from the time dependence of these oscillations,\textsuperscript{5.4-5.9} an independent measurement of the density from the period of the oscillations,\textsuperscript{5.10} and the best measurement of the renormalized electron mass.\textsuperscript{5.11}

The usual de Haas type oscillations in a metal are ascribed to density of states changes at the Fermi level as the Landau levels cross through the Fermi level with no change in actual density of the electron gas in the metal. The case for the electron-hole liquid is somewhat different. As the magnetic field changes the free energy must be recalculated and minimized to obtain the appropriate density and energy for the liquid. Theoretically, this has only been done for low fields,\textsuperscript{5.3} $\hbar \omega_c \ll E_F$, where $E_F$ is the Fermi level, and

$$\hbar \omega_c = eH/m_c$$

(V.1)

is the cyclotron frequency with $m_c$ the cyclotron effective mass; and very high fields,\textsuperscript{5.12} $\hbar \omega_c \gg E_F$. Experimentally, however, the
situation is much clearer. It has been found that the Fermi level of the EHL remains quite constant in a magnetic field, and the density oscillates periodically with period $E_F^e / \hbar \omega_c \sqrt{V}$, where $E_F^e$ is the electron Fermi level within the EHL. Effects due to the holes do not enter until $H > 30 \text{kOe}$ because structure for the holes is found only at these higher fields. At lower fields the holes only provide a background change due to the complicated structure of the hole Landau levels.

As a consequence of these density oscillations, the luminescence intensity oscillates because the radiative recombination rate

$$\tau_r^{-1} = Bn,$$  \hspace{1cm} (V-2)

where $n$ is the density, and $B$ is the radiative recombination coefficient.

Because in the steady state the luminescence intensity

$$I_L = N \tau_r^{-1},$$  \hspace{1cm} (V-3)

and the total number of pairs in the EHL, $N$ is proportional to $\tau_0$, the total recombination lifetime for the EHL, the situation is slightly more complicated for this case. The total lifetime can be written

$$\tau_0 = (A + Bn + Cn^a)^{-1}$$  \hspace{1cm} (V-4)

where the density independent recombination rate is defined

$$A = \tau_i^{-1}$$  \hspace{1cm} (V-5)

where $\tau_i$ is the density independent lifetime; and the Auger recombination rate is defined

$$Cn^a = \tau_a^{-1}$$  \hspace{1cm} (V-6)

where $\tau_a$ is the Auger lifetime, $C$ the Auger coefficient and $a$ the appropriate power for the dominant Auger process. This process is
believed to be a phonon assisted three-body process for the EHL in GeV.14 yielding \( a = 2 \) for this case. Efficiencies for these processes can be defined as follows:

\[
\begin{align*}
\eta_i(n) & = \frac{\tau_o(n)}{\tau_i(n)} \\
\eta_r(n) & = \frac{\tau_o(n)}{\tau_r(n)} \tag{V-7} \\
\eta_a(n) & = \frac{\tau_o(n)}{\tau_r(n)}.
\end{align*}
\]

This leads to a steady state luminescence intensity

\[
I_{ss} = \eta_r(n) = \frac{Bn}{(A + Bn + Cn^a)^{-1}} \tag{V-8}
\]

Thus for any EHL with \( \eta_r < 1 \), there will be a change in steady state luminescence intensity with a change in density. In a magnetic field this implies that \( I_{ss} \) will be periodic with period \( E_F^e/\hbar \omega_c \).

By measuring the period of \( I_{ss}(H) \) in \( 1/H \) it is possible to measure \( E_F^e \) quite accurately. Because \( E_F^e \) is independent of \( H \), and

\[
E_F^e = \frac{\hbar^2}{2m_e^*} \left( \frac{3\pi^2 n_o}{2} \right)^{2/3} \tag{V-9}
\]

where \( m_e^* \) is the density of states effective mass for electrons, it is possible to get an accurate value for \( n_o \), the equilibrium density at \( H = 0 \). Therefore we measured the luminescence intensity of the SCEHL as a function of \( H \). The experimental set up was the same as described in Chapter III for imaging experiments with the following small changes. No imaging or spatial resolution was used, and in order to maximize the sensitivity of the system to changes in the lifetime, the lock-in was set up differently. Usually a lock-in phase is set in order to maximize the signal. In this case, however, since our lock-in (PAR model 186 or 128A), is sensitive to the shape of the incoming signal and since the lifetime of the EHL (~500 \( \mu s \))
is an appreciable fraction of the on or off time (~2 ms) of the chopped laser excitation, it was found that the best sensitivity to the oscillations was obtained by setting the phase to null the signal at H = 0. This occurs because the oscillations with H are due to changes in lifetime which translate to changes in shape for the lock-in. Other changes due to the magnetic field, such as change in pumping efficiency, which do not cause a change in shape of the signal therefore do not contribute directly to the signal. This allowed the use of low absorbed laser power $P_a$.

Other magnetic field scans were taken by replacing the lock-in with a boxcar integrator which I designed and built (Fig. V-1). This boxcar integrator is based on the gated integrator and digital delay modules produced by Evans Associates. The delay module positions the gate very accurately ($\pm 10$ ns) by using a 10 MHz clock and special interpolation circuitry. Two gated integrator modules are used with a common preamp to give automatic baseline corrections. It was found that the hold time and leakage of these modules was as good as or better than the PAR 160 boxcar. Special care was taken to give the preamp high gain (up to $\times1000$), maximum gain flexibility ($\times2$, $\times5$, and $\times10$), and good frequency response (DC $\rightarrow$ 10 MHz for gain $< 100$, $\sim 1$ MHz for gain $> 100$). A very fast trigger ($\sim 2$ ns) enables the minimum delay to be quite short ($\sim 500$ ns). This minimum delay is independent of the delay time on the digital channel (A channel). The B channel has an analog delay governed by an R/C time constant set with front panel switches. This channel is designed to monitor the baseline, and the usual output of the boxcar is A-B. The gain of the boxcar is entirely in the preamp. The time delay of A is displayed digitally on the front panel. The output is displayed by using a digital voltmeter on the front panel. Sweeping and external control are made possible through rear panel connections.
Magnetic field scanning and data collection was made more automatic with the ramp digitizer which I designed and built. The current monitor output of the magnet power supply is digitized and displayed by a digital voltmeter. The digital result is compared with a digital reference; when the two are equal a pulse is generated, and the reference is changed by the amount set on the front panel. In this way the ramp digitizer generator pulses in equal magnetic field increments.

Both the lock-in and boxcar techniques yield oscillations with the same period and phase. Figure V-3 shows three sets of oscillations for a sample stressed along $\langle 111 \rangle$, for different orientations of the magnetic field with respect to the crystal axes. Trace (a) $\hat{H} \parallel \langle 001 \rangle$, and trace (b) $\hat{H} \parallel \langle 110 \rangle$, were taken for $P_a = 3.2 \text{ mW}$, using the lock-in technique described above. Trace (c) $\hat{H} \parallel \langle 111 \rangle$, was taken for $P_a = 2.4 \text{ mW}$ at a time $t = 800 \mu s$ after the light was shut off using the technique with the boxcar. The period of the oscillations changed consistently with the angular variation in the electron cyclotron mass, assuming the electron masses are unchanged from bulk Ge. I fit these curves using a magneto-oscillatory expression similar to that reported by Keldysh and Silin.\textsuperscript{V.1} I assumed that only one electron valley was occupied and that the oscillations were due to the electrons and not the holes as discussed above. From these fits I find the electron Fermi level and average value for the $e$-$h$ pair density to be

$$E_F^e = 2.3 \pm 0.12 \text{ meV}$$

$$n_0 = (5.2 \pm 0.5) \times 10^{16} \text{ cm}^{-3} \quad \text{(V-10)}$$

$$T = 1.6 \text{ K}, -\gamma_m \approx 6 \text{ kfg/mm}^2$$
where $\sigma_M$ is the applied stress.

The above result was obtained at a sufficiently low excitation level that the luminescence linewidth (Ref. V.15, Fig. 10) was not broadened, so the above density is considered to be the equilibrium value. At higher excitation levels the period of the oscillations was observed to increase with $P_a$, indicating compression of the liquid. Figure V-4 shows a magnetic field sweep for the same field orientation and experimental technique as in Fig. V-3c but taken at a much higher excitation level ($P_a = 120$ mW, $t = 100$ µs). For this data we found $E_F^e = 3.68$ meV, corresponding to $n = 1.05 \times 10^{17}$ cm$^{-3}$. It is perhaps surprising that, with such a wide range of density within the SCEHL, the oscillations are still very well resolved. While this represents some kind of an average density, it shows clearly that significant compression can easily be obtained with moderate excitation levels, as discussed in Chapter IV.

As was discussed in Chapter I, it is clear theoretically that only one conduction band valley will be occupied in the strain well. This can be tested indirectly by comparing the results of the luminescence lineshape fits in Chapter IV with the results of Eq. (V-10). It can be seen clearly that because the magneto-oscillatory luminescence data give a higher value for the electron Fermi level, $E_F^e = 2.3$ meV, than the spectroscopic data, $E_F^e = 2.1$ meV, which assumed only one electron valley occupied, no more than one valley can be occupied. This conclusion must be considered somewhat indirect, however, due to the different temperatures and possibly different strains used in the two experiments. I therefore decided to check the degeneracy of the valleys directly by measuring the angle dependence of the magneto-oscillations.
In order to more carefully check this angle dependence, I designed the special sample rotation head pictured in Fig. V-5. Angle dependent data of extrema positions in field are shown in Fig. V-6. The solid curves are the expected extrema positions assuming an unchanged cyclotron mass. An extrema will occur wherever

\[ 2(j + 1)\hbar \omega_c = E_F^e \]  

(V-11)

where \( j \) is an integer. Therefore the field at a given extrema \( H_e \propto m_c \) from Eq. (V-1), and the fact \( E_F^e \) is constant. However, this is not conclusive. Figure V-7 shows the period in units of \( E_F^e m_c/m_o \) obtained from fitting the data to magneto-oscillatory function as above. It can be seen that although there is no significant deviation from the expected behavior, there seems to be a trend toward \( E_F^e m_c \) being greater than expected at larger angles. Since the expected deviations due to the renormalized mass would be insignificant at these angles, it is most likely that one of two things is happening. First, it is possible that the density is changing due to pumping efficiency changes with angle, or due to magnetostriction effects (see Chapter III and references III.10 and III.11). Second, the assumption that \( E_F^e \) is independent of angle may be breaking down. In order to check this assumption, it is necessary to recalculate the free energy as a function of magnetic field and density and then minimize it. This theoretical task is beyond the scope of this thesis.

Due to the above effects it is not possible to determine the renormalized electron mass parameters in the SCEHL. However, it is clear from the angle dependence that only one electron valley is occupied. There are no observable oscillations due to electron valleys other than the valley selected by the (111) stress as predicted by Markiewicz et al. V.16
In order to obtain information about the decay processes of the SCEHL magneto-oscillation experiments were performed at discrete times after a 100 ns pulse from a GaAs laser. Under these conditions the drop grows very rapidly to a maximum size as seen in Chapter II and then decays in the same way as a drop formed by CW excitation. The total energy per pulse was chosen to produce a drop with radius \( R \approx 125 \mu \text{m} \). The magneto-oscillations of the luminescence at several delay times are shown in Fig. V-8 with \( H \lesssim 18 \text{kOe} \) and \( H \parallel (1\overline{1}0) \). It can be seen that at \( t = 1 \text{ ms} \) the oscillations are reduced in amplitude but have not changed sign, where \( t = 1 \text{ ms} \) is approximately twice \( \tau_0(n_0) \). This is in contrast with the case for unstressed Ge, where the oscillations change sign after approximately one-half of the zero field decay time.\(^4\)

To understand the significance of this result, it is necessary to consider the kinetics of a pulsed experiment. Because the drop lifetime is much longer than the laser pulse width, \( \approx 100 \text{ ns} \), the drop initially contains a constant number \( N(t=0) \) of e-h pairs, assuming the pumping efficiency is independent of magnetic field. The luminescence intensity as a function of time is given by

\[
I(H,t) = N(t,n) = BnN(0)e^{-t/\tau(n)}. \tag{V-12}
\]

At a magnetic field for which the equilibrium density is higher than the zero field value \( n_0 \), the initial luminescence intensity is greater than \( I(0,0) \), and the lifetime \( \tau(n) \) is shorter than \( \tau_0(n_0) \). Therefore at some later time \( t_t \),

\[
I(0,t_t) = I(H,t_t); \tag{V-13}
\]

at this turnover time the magneto-oscillations change sign. Assuming \( n(H) - n_0 \ll n_0 \) a straightforward analysis yields

\[
\tau_t^{-1} = Bn_0 + aCn_0^3. \tag{V-14}
\]

This result combined with Eqs. (V-4) and (V-7) yields
\[ \tau_o(n_0)/\tau_t = \xi + a\xi (n_0) \equiv \xi. \quad (V-15) \]

The quantities \( \xi \) and
\[ 1 - \xi = \xi_1 - (a - 1)a\xi (n_0) \quad (V-16) \]
appear throughout the analysis of the time behavior of magneto-
oscillatory phenomena. Other combinations of recombination effi-
ciencies cannot easily be obtained from these experiments. If
\( \xi < 1 \), it can be seen from Eq. (V-16) that the density independent
recombination is more important than the Auger mechanism. Conversely,
if \( \xi > 1.0 \), the Auger mechanism is more important than the density
independent mechanism (for \( a = 2 \)). The data displayed in Fig. V-8
give direct information about \( \xi \). Since the turnover time \( \tau_t \geq 2\tau_o(n_0) \)
we conclude that \( \xi \leq \frac{1}{2} \). In addition, from Eq. (V-16) I find that
\[ \xi_i \geq 0.5, \quad (V-17) \]
in agreement with lifetime studies in very large compressed volumes
of SCEHL.\textsuperscript{17} We find from the data of Betzler et al.\textsuperscript{4} that \( \xi \sim 2 \)
for the EHL in unstressed Ge. This indicates that the Auger mechanism
dominate for this case.

Betzler et al.\textsuperscript{4} concluded that the density independent recom-
bination efficiency was less than 10\% of the Auger efficiency and
therefore set \( \xi_i = 0 \). However, we note that if \( \xi_i = 0.05 \) and \( \tau_o = 36 \mu s \) for unstressed Ge, this corresponds to a value of \( \xi_i = 720 \mu s \).
For the SCEHL, the same density independent decay time corresponds
to \( \xi_i = 0.69 \) (SCEHL) using \( \tau_o(n_0) = 500 \mu s \). It is clear that such
a process which can be neglected in unstressed Ge can at the same
time dominate the recombination for the SCEHL.

A density independent recombination time of several hundred \( \mu s \)
is reasonable for Ge. Possible recombination sites include shallow
and deep traps. Several groups have measured the EHL decay time in doped (unstressed) Ge. While there is some variation in the results, the lifetime generally starts to decrease when the impurity concentration, \( N_i \), is greater than \( \sim 10^{15} \) - \( 10^{16} \) cm\(^{-3}\). For example, Zhurkin et al. report \( \tau_0 = 24 \) ms for \( N_i = 2 \times 10^{16} \) cm\(^{-3}\).

As impurities. Assuming that the e-h pair density remains constant at these doping levels, and that \( \tau_i \propto N_i \), the decrease in lifetime may be attributed to impurity-induced recombination. Therefore \( \tau_i \sim 70 \) ms for \( N_i = 2 \times 10^{16} \) cm\(^{-3}\) implies \( \tau_i \sim 700 \) ms corresponds to \( N_i \sim 2 \times 10^{15} \) of shallow impurities. Because this is at least three orders of magnitude larger than the observed impurity densities in this crystal, shallow impurities cannot explain our data.

On the other hand, an estimate can also be made for deep levels, where nonradiative recombination could take place via multiphonon emission. Experimental values for nonradiative capture cross-sections at room temperature have been tabulated for Ge and range from \( \sigma \sim 10^{-14} \) - \( 10^{-16} \). A simple model predicts that the capture rate is \( \sim vN_i \), where \( v \) is the carrier velocity which may be taken to be the Fermi velocity for carriers in the EHL. In this case \( \tau_i \sim 1 \) ms would lead to \( N_i \sim 10^{10} \) cm\(^{-3}\). The samples have not been characterized to this level, although it is known that levels are present associated with hydrogen. Thus the study of the SCEHL in differently prepared samples may prove to be a sensitive test of impurity-induced recombination and provides a good test of the importance of this mechanism for the EHL.
CHAPTER V: References


V.18 C. Benoit à la Guillaume, and M. Voos, Solid State Commun. 11, 1585 (1972).


V.20 E. E. Haller, private communication.

CHAPTER V: Figure Captions

Fig. V-1 Block diagram of digital boxcar which does automatic baseline subtraction (B channel).

Fig. V-2 Block diagram of the ramp digitizer which is used to put marks on the chart recorder and control other external logic. It is used in conjunction with an analog ramp generator used to control the superconducting magnet power supply.

Fig. V-3 Total luminescence intensity as a function of magnetic field for three different orientations of the field with respect to the crystal axes. (a) $\mathbf{H} \parallel \langle 001 \rangle$, (b) $\mathbf{H} \parallel \langle 110 \rangle$, for $P_a = 3.2 \text{ mW}$, at steady state (c) $\mathbf{H} \parallel \langle 111 \rangle$, for $P_a = 2.4 \text{ mW}$ and $t = 800 \mu\text{sec}$. Permanent stress geometry (stress along $\langle 111 \rangle$). $T = 1.6 \text{ K}$. The vertical scale is offset from zero by an arbitrary amount. The oscillations are about a 10\% effect. Sample CR36.

Fig. V-4 Total luminescence intensity as a function of magnetic field for the same conditions as for Fig. V-3 (c), except that $P_a = 120 \text{ mW}$, $t = 100 \mu\text{sec}$. The peaks have moved to higher fields, showing compression of the liquid. The vertical scale is offset from zero by an arbitrary amount.

Fig. V-5 Photograph of the special sample rotater. The sample is visible in the lower portion of the lefthand view. Pumping light is directed onto the sample by the $45^\circ$ mirror in the middle of the rotater (righthand view). Luminescence is collected by the $45^\circ$ mirror at the bottom of the rotater.

Fig. V-6 Extrema in the magneto-oscillations plotted as a function of angle measured from the $\langle 111 \rangle$ axis. $P_a = .32 \text{ mW}$. Permanent stress geometry (stress along $\langle 111 \rangle$). Solid curves
are the expected angle dependences for these extrema assuming no mass renormalization.

Fig. V-7 Period of the magneto-oscillations as a function of angle. Solid line is an approximate fit to the data with $E_F^0 = 2.54$ meV. No mass renormalization was used. Experimental conditions are for the previous figure. Sample CR36.

Fig. V-8 Total luminescence intensity as a function of magnetic field $H$ for several different delay times after a 100 ns excitation pulse, absorbed energy $\sim 1.5$ ergs. $H \parallel (\overline{110})$; $T = 1.6$ K. The vertical scale is offset from zero by an arbitrary amount. The symbols at the right indicate 10% of the total intensity. Sample CR50.
Fig. V-1
Fig. V-2
Fig. V-6

EXTREME POSITION $H_e$ (kOe)

ANGLE FROM $\langle 111 \rangle$ (deg)
Fig. V-7
Fig. V-8

XBL 7812-13245
CHAPTER VI

The Metal-Insulator Transition in Uniformly Stressed Ge
Metal-insulator (MI) transitions have been of constant interest since they were first proposed by Mott. Experimental studies of the FE-EHL phase diagram in Ge have received attention due to the interest in what aspects, if any, of a Mott M-I transition are also present in the data. In particular, the nature of the gas phase near the critical point for EHL formation has been questioned. Rice has pointed out that there may be two critical points for the EHL system as was first proposed by Landau and Zeldovich for liquid Hg. Various static screening models for the expected ionization catastrophe have been proposed based on Debye-Hückel static screening. These give a metal-insulator transition density of

\[ n_m = \frac{k_B T}{E_{bo}} \frac{1}{11\pi a_{bo}^3} \]  

where \( E_{bo} \) and \( a_{bo} \) are exciton binding energy and Bohr radius at zero density, respectively.

Experimental results, for instance those of Thomas and Rice in Ge[4,2] and Shah et al. which do not agree with these simple static screening models have lead to more sophisticated models which include important dynamic screening effects and self energy corrections. (See Ref. and the citations therein.) The latest model by Zimmerman et al. finally seems to explain the most striking experimental result that the absolute energy of the exciton as observed in luminescence or absorption does not change as the exciton density approaches \( n_m \) from Eq. (VI.1). The results of reference show that the exciton energy does not depend on the density of dissociated pairs, \( n_f \) because the exciton is a neutral species and does not couple strongly to the charge density oscillations
of the free pairs. However, the energy per pair of these disassociated pairs, $E_f$, depends markedly on $n_f$ due to self energy corrections.

Taking the functional form from figures 1 and 2 of Reference VI.8 and $n_m$ from Eq. (VI-1), I find I can approximate these results for $E_f(n_f)$ with the simple expression

$$E_f = E_{bo}[1 - (n_f/n_m)^{1/2}], \quad \text{(VI-2)}$$

where the zero of energy is taken to be the exciton energy.

As was first pointed out by Balslev, VI.10 this type of density dependent energy can lead to interesting thermodynamic results.

Assuming nondegenerate species, the free energy for a system of free pairs and excitons is

$$F = N k_B T((1 - \eta)[\ln(V_{Q_x} n_t (1 - \eta)) - 1]$$

$$+ \eta[\frac{E_{bo}}{k_B T}1 - (\frac{n_f}{n_m})^{1/2}] + \ln(V_{Q_e} V_{Q_h} n_t^2 n_x^2) - 2}] \quad \text{(VI-3)}$$

where $N$ is the total number of pairs, $T$ is the temperature,

$$\eta = n_f/n_x \quad \text{(VI-4)}$$

is the ionization ratio of the density of free pairs, $n_f$, to the density of excitons, $n_x$. $n_m$ is from Eq. (VI-1), $n_t = n_f + n_x$ and

$$V_{Qa} = \frac{1}{\nu_a} \left(\frac{2 m a^2}{m a k_B T}\right)^{3/2} \quad \text{(VI-5)}$$

is the quantum volume of species $a$. $\nu_a$ and $m_a$ are degeneracy and density of states mass of species $a$. Excitons, electrons, and holes are referred to by $x$, $e$, and $h$, respectively. By minimizing this free energy with respect to $\eta$, or equivalently setting the chemical potential of the free pairs and excitons equal one obtains the following implicit equation for $\eta(n_t)$.
\[ 1 - \eta - n_t \eta^2 \frac{V_x V_q}{V_{\text{Qx}}} \exp\left(\frac{E_{\text{bo}}}{k_B T} \left[1 - \frac{3(n_t \eta)}{2(n_m)}\right]\right) = 0. \] (VI-6)

In Fig. VI-1 the solutions of this equation \( n(n_t) \) are plotted for various temperatures for excitons in Ge[1,1]. It can be seen that at low temperatures there is a bistability region where there are three solutions: two stable solutions and one unstable solution. At higher temperatures there is only one solution and the bistability region is gone. The absolute limits of this bistability region are defined by the singularities in \( \frac{dn}{d n_t} \). This bistability region is plotted in Fig. VI-2 along with the gas liquid phase diagram for the same conditions from Feldman et al. VI.12 It can easily be seen that a considerable bistability region may be observable above the critical temperature for EHL formation. The actual limits of the bistability region, region 3, would have to be determined by nucleation theory and include an interface energy between phase 1, the insulating phase which is predominantly excitons, and phase 2, the metallic phase which is predominantly free pairs. Further, the effects of EHL condensation and the liquid-gas coexistence region, region 4, have to be taken into account. However, it is possible to get an estimate of the stability of phases 1 and 2 by examining the free energy. The average free energy per pair, \( F/N \), is plotted versus \( \eta \) for various densities in Fig. VI-3. It can be seen that, unless there are substantial changes to this free energy due to the corrections mentioned above, there is almost no barrier to formation of the metallic phase from the insulating one once there are two allowed phases. The metallic phase is energetically favorable once \( n_t \) is somewhat greater than \( n_m \). Further this phase seems to be much more stable, being quite bound with respect to the insulating phase well into the bistability region. Thus the bistability region
would be modified to a small coexistence region near the boundary between regions 1 and 3, and the boundary between regions 2 and 3 would be moved much closer to $n_m$.

In order to gain more experimental information about the metal-insulator transition in semiconductors, Dr. I. Balslev and I have studied the luminescence from highly excited germanium under a uniform $(111)$-uniaxial stress of $16 \text{ kgf/mm}^2$. As was discussed in Chapter 1, this stress changes the electron-hole configuration to Ge[1,1]. As a result the EHL phase diagram in unstressed Ge, with a liquid density of $n(T=0) = 2.3 \times 10^{17} \text{ cm}^{-3}$ and critical temperature of $T_c = 6.5 \text{ K}$ is shifted to markedly lower densities and temperatures: $2 \times 10^{16} \text{ cm}^{-3}$ and $3.5 \text{ K}$. Furthermore, stress changes the values of $E_{bo}$ and $a_{bo}$, yielding M-I transition densities from Eq. (VI-1):

<table>
<thead>
<tr>
<th></th>
<th>Ge[4,2]</th>
<th>Ge[1,1]</th>
</tr>
</thead>
<tbody>
<tr>
<td>$E_{bo}$</td>
<td>4.1 meV</td>
<td>2.65 meV</td>
</tr>
<tr>
<td>$a_{bo}$</td>
<td>112 Å</td>
<td>176 Å</td>
</tr>
<tr>
<td>$n_m(4K)$</td>
<td>$1.72 \times 10^{15} \text{ cm}^{-3}$</td>
<td>$6.86 \times 10^{14} \text{ cm}^{-3}$</td>
</tr>
</tbody>
</table>

Stressing Ge lowers $n_m$ by a factor of 3 and allows us to more readily reach the M-I transition density. Since when $T < T_c$ the EHL luminescence dominates the observed spectrum, it is very helpful to take data above $T_c$ to observe the M-I transition. Because $T_c < 4.2 \text{ K}$, the present experiments can be done immersed in liquid helium rather than gas allowing better temperature control at high excitation levels.

We chose to use a long rectangular ultrapure Ge sample of dimensions $0.7\times0.7\times10 \text{ mm}^3$ with the long dimension along a $(111)$ axis so that we would have very uniform stress near the center of the sample due to the large aspect ratio of greater than 10:1. Larger samples
with a similar aspect ratio did not seem to work as well, possibly due to sample preparation problems. Ordinary tin solder was used as a pad between the stressing rods and the sample because this material is plastic at room temperature conforming to the surface of the Ge. At low temperature it becomes quite hard and transfers the force to the Ge sample without further deformation. The stress was applied with special precision stressing apparatus I designed, pictured in Fig. VI-4. This head will apply very precise variable uniaxial forces to above 100 kgf. By using appropriate attachments both uniform and nonuniform stresses may be easily applied.

In order to insure stress reproducibility between runs and to cancel any effect of different helium vapor pressures, a birefringence sampling apparatus was added to the experimental setup shown in Fig. VI-5. By checking the birefringence, the stress could be reproduced to better than ±0.05 kgf/mm². The stress of 16 kgf/mm² was chosen to correspond to the stress used by Feldman et al. VI.12 in measuring the EHL phase diagram in Ge[1,1] allowing us to use this phase diagram to calibrate our exciton densities. For our sample this stress also corresponds to the third half integer retardation between the ordinary and extraordinary wave at \( \hbar \omega = 700 \text{ meV} \). This minimum in transmitted light intensity is a very convenient place to obtain reproducible birefringence measurements. The sample was pumped with a chopped CW 3W argon ion laser focused to a spot of less than 0.5mm in diameter onto the central region of the sample. Two-thirds of the openings on the chopping wheel were covered to give a 0.7 ms pulse and a 17% duty cycle. This was done to minimize the sample heating due to the high laser powers needed to obtain M-I transition densities. The collected luminescence was spectrally analyzed using
the stepper motor-driven spectrometer and collected in a multichannel analyzer as pictured in Fig. VI-5 and described in detail by Westervelt in his thesis.

Data obtained for 0.07 W absorbed in the sample are shown in Fig. VI-6. Dots correspond to data points, and the smooth curves are exciton lines which approximately fit the data. A triangular slit function with a halfwidth of 0.4 meV was used to account for the 150μ slits in the experiment. A further 0.6 meV of Gaussian broadening was also added to give consistent results for temperatures and line positions. At 0.07 W absorbed the temperature of the exciton gas was found to be essentially the same as the He bath temperature. At 2.19 K the EHL line at 700 meV is also observed (Fig. VI-6a); however, at 3.15 K (Fig. VI-6b) this line is very weak, and at 4.2 K (Fig. VI-6c) is nonexistent. This is the expected behavior for the EHL phase with a critical temperature of 3.5 K.

Data for higher absorbed power, 0.28 W, are shown in Fig. VI-7, with dots and smooth curves as for Fig. VI-6. At this absorbed power considerable beating, 2 to 3 K, is evident in the exciton gas. Further, a second line, which we interpret as being due to the high density phase of a M-I transition, is evident at about 701 meV. This line shows very little temperature dependence.

We estimate the average total pair density $n_t$ (i.e., free carriers, excitons, EHL) due to the exciting light in two independent ways. First, using the absorbed light power $P_a = 0.27$ W ($h\omega = 2.4$ eV) with a conservatively estimated carrier lifetime of 2 μs, and an effective volume of 0.7 mm$^3$ we arrive at a density of $2 \times 10^{15}$ cm$^{-3}$ assuming 100% production efficiency. We obtain the second estimate by using the phase diagram of Feldman et al. EHL is formed at 2.19 K and
\( P_a = 0.07 \text{ W}, \text{ implying a density greater than } 1.6 \times 10^{15} \text{ at } P_a = 0.27 \text{ W.} \)

We interpret the luminescence line at 701 meV shown in Fig. VI-7 as being due to free pairs by a process of elimination. The line cannot be due to EHL since it has the wrong temperature dependence and exists above the critical temperature for the EHL. It is very unlikely that the line is due to larger complexes such as biexcitons as was proposed by Thomas et al. VI.6 Biexcitons have the largest binding energy of these complexes, but this energy is still approximately equal to the temperature VI.16, and they should be completely disassociated.

Because the line at 701 meV is spectrally resolved and appears only at high pump powers, we interpret the line as due to the high density phase of a M-I transition. Due to the inhomogeneous nature of the experiment it is impossible to fit this line. However, the shape does give good qualitative evidence that our interpretation is correct. The line is broad with a pronounced low energy tail. This shape can be understood qualitatively by assuming a nonuniform plasma: the emission from dense regions is broad and at low energies (large band renormalization), while emission from less dense regions is narrower and at higher energies.
CHAPTER VI: References


VI.2 N.F. Mott, Metal Insulator Transitions (Barnes & Noble, New York, 1974).


~ 3 μs, is probably due to surface recombination presumably of excitons. Therefore our estimate of 2 μs is quite reasonable considering the small sample used in this case.

VI.16 Reference VI.6 indicates a binding of 2 K in unstressed Ge.

This should be even smaller in stressed Ge.
CHAPTER VI: Figure Captions

VI-1 Predicted ionization ratio \( n \), defined in the text versus total density \( n_t \) for Ge[1,1], \( m_x = 0.419 \, m_o \), \( m_e = 0.218 \, m_o \), \( m_h = 0.088 \, m_o \), \( n_x = 4 \), \( n_e = 2 \), \( n_n = 2 \), \( E_{bo} = 2.65 \, \text{meV} \), \( a_{bo} = 176 \, \text{Å} \). Curve:
   a) \( T = 3 \, \text{K} \); b) \( T = 5 \, \text{K} \); c) \( T = 7 \, \text{K} \).

VI-2 Phase diagrams for the metal insulator transition from Eq. (VI-6.)
The dotted line is \( n_m(T) \) from Eq. (VI-1). Region 1 is the insulating phase predominantly excitons, Region 2 is the metallic phase predominantly free pairs, Region 3 is the maximum extent of the bistability region where phase 1 or phase 2 is stable depending on past history. This overlaps region 4 which is the gas, electron-hole liquid coexistence region taken from reference VI.12.

VI-3 Average free energy per pair for Ge[1,1] at 3 K as a function of ionization fraction \( n \) for various densities; a) \( n_t = 10^{14} \, \text{cm}^{-3} \), b) \( n_t = 4 \times 10^{14} \, \text{cm}^{-3} \) showing the low density limit of the bistability region, c) \( n_t = n_m = 5.1 \times 10^{14} \, \text{cm}^{-3} \), d) \( n_t = 10^{15} \, \text{cm}^{-3} \). The symbol + indicates a maximum in the free energy and the symbol - indicates a minimum in the free energy.

VI-4 Assembly drawing of precision stressing head.

VI-5 Schematic of apparatus used in these experiments. For checking the birefringence the 45° mirror and crossed polarizers were inserted, and the laser was blocked. For taking spectra the 45° mirror and the crossed polarizers were removed, and the incandescent lamp was turned off.

VI-6 Luminescence spectra with \( P_a = 0.07 \, \text{W} \). The dots represent data points and the smooth curves are exciton lines which approximately fit the data with a slit width of 0.4 meV and 0.6 meV of Gaussian
broadening. The large tick at 701.8 meV shows the low energy edge of all the exciton lines. a) He bath at 2.19 K; exciton lineshape at 2.5 K. Notice characteristic EHD luminescence at 700 meV. b) He bath at 3.15 K; exciton lineshape at 3 K. c) He bath at 4.25 K; exciton lineshape at 4 K.

VI-7 Luminescence spectra with $P_a = 0.28$ W. All symbols and parameters as in Fig. VI-6 except temperature. a) He bath at 2.1 K; exciton line shape at 4 K. b) He bath at 3.15 K; exciton line shape at 6 K. c) He bath at 4.25 K; exciton lineshape at 7 K. Notice the temperature independent resolved peak at 701 meV which is interpreted as being due to the high density metallic phase of a M-I transition.
Fig. VI-1
Fig. VI-2
Fig. VI-3
Fig. VI-5
Fig. VI-6
Fig. VI-7

(c)  
BATH: 4.26 K  
FIT: 7 K

(b)  
BATH: 3.12 K  
FIT: 6 K

(a)  
BATH: 2.19 K  
FIT: 4 K

ENERGY (meV)

LUMINESCENCE INTENSITY (Rel. units)
CHAPTER VII

Conclusions
A number of experiments have been presented here which illustrate the nature of the excited metallic phases in strained Ge. First, the formation processes for the strain-confined electron-hole liquid was studied. It was found that the photo-excited pairs formed at the crystal surface move quickly into the strain potential well formed by the Herzian contact stress. Both the phonon wind from the absorbed laser light and the strain potential gradients are important for this motion. The mass anisotropy of the electrons was found to have a profound effect on the motion of material into the SCEHL, focusing it along (111) strain axes.

Next, the remarkable magnetostriction of the SCEHL was discovered and thoroughly explained. The recombination currents which are necessary to maintain the EHL at constant density are deflected by the magnetic field such that resulting magnetization is paramagnetic. The SCEHL then deforms to lower its total energy by increasing its size perpendicular to the field. The details of the effect were shown to agree very well with Markiewicz's theory and show conclusively that this magnetostriction is due to this dynamic recombination magnetization.

Studies of the luminescence from the SCEHL established its density, $n_0 = 4.6 \pm 0.3 \times 10^{16}$ cm$^{-3}$, and condensation energy, $\phi = 10.8 \pm 1$ K at $T = 3.3$ K and $\sigma_m \approx -6$ kgf/mm$^2$, by fitting the lineshapes of the FE and SCEHL. An independent measure of the condensation energy, $\phi = 9.2 \pm 1$ K, in agreement with the spectroscopic value was found by studying the threshold for EHL formation as a function of temperature.

More magnetic studies on the SCEHL to study the magneto-oscillatory luminescences provided an independent measure for the density of the
SCEHL, \( n_0 = 5.2 \pm 0.5 \times 10^{16} \text{ cm}^{-3} \) at \( T = 1.6 \text{ K} \) and \( \sigma_m \approx -6 \text{ kgf/mm}^2 \) in agreement with the spectroscopic determination. By studying the angle dependence of these oscillations the fact that only one electron valley is occupied in the SCEHL was conclusively demonstrated. Finally, by studying the time dependence of these oscillations it was discovered that a density independent recombination mechanism is important in the SCEHL decay.

By studying highly excited uniformly strained Ge, a possible metal-insulator transition was observed above the critical point of the EHL. Because a spectroscopically well-resolved approximately temperature independent line was observed at high pump powers, it is believed that this M-I transition was observed. The data is compared with the latest theories, and the agreement is found to be quite good.

Further work is indicated in a number of areas. A more thorough study of the phase diagram of the SCEHL is indicated to include a stress dependence of \( n_0 \), \( \phi \), and the critical temperature \( T_c \). More studies of the magneto-oscillatory luminescence could possibly result in the determination of the renormalized carrier masses in the SCEHL. Finally, a fair amount of experimental and theoretical work on the metal-insulator transition is indicated including studies in the strain well.
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